

# Parity Violating Effects in Elastic Electron Deuteron Scattering

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The general expressions for parity violation observables in elastic scattering of polarized and/or unpolarized electrons from unpolarized deuterons are given and are numerically evaluated for the kinematics of SAMPLE, PVA4 and G0 experiments. The dominant contribution from the interference of  $\gamma$  and Z exchange as well as the smaller contributions from strangeness ( $s\bar{s}$ ) components of the nucleon, parity odd admixtures in the deuteron wave function, anapole moments and radiative corrections are included and discussed in the context of parity violating electron scattering experiments of present interest.

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## I. INTRODUCTION

Parity violating electron scattering of polarized electrons from nucleons and nuclei provides an important tool to probe the electroweak structure of hadrons. The dominant contribution to the parity violating observables in electron scattering processes comes from the interference of the electromagnetic amplitude given by one photon exchange and the weak neutral current amplitude given by Z exchange. These processes are therefore considered to be ideal for studying the structure of weak neutral currents [1, 2]. A knowledge of neutral weak currents is of special importance in understanding the role of sea quarks in the structure of nucleons [3]. The isoscalar sector of the weak neutral current contains information about the strangeness ( $s\bar{s}$ ) content of the nucleon which is manifest through the isoscalar vector and axial vector form factors  $G_E^S$ ,  $G_M^S$ , and  $G_A^S$ . A theoretical and experimental study of these form factors provides thus information about the role of strange sea quarks in the nucleon. The initial evidence for the nonzero strangeness ( $s\bar{s}$ ) content of the nucleon as seen in deep inelastic scattering (DIS) [4], neutrino scattering [5] and pion nucleon scattering [6] has generated great interest in looking for these effects through parity violating observables in elastic and inelastic scattering of electrons from nucleons and nuclei.

Many experimental programs have been started at various electron accelerators to search for the strangeness content of the nucleon through the observation of parity violating asymmetries in the scattering of polarized electrons from proton [7]-[14], deuteron [8], [15, 16] and  ${}^4\text{He}$  [11, 16] targets from which first results have been reported for the strange form factors [17]-[19]. The initial results for the electron asymmetry in quasi-elastic scattering of polarized electrons from deuterium targets in the SAMPLE II experiment indicated that the effect of radiative corrections in the weak axial coupling constant  $G_A^S$  could be large, making it difficult to extract information on the strangeness charge and magnetic moment form factors  $G_E^S$  and  $G_M^S$ . Later, the updated results from SAMPLE II and the new results from SAMPLE III [15, 16] have shown that the contribution to the asymmetry from radiative corrections [19]-[23] to the weak coupling constants, especially the contributions like anapole moments [21]-[23] of the deuteron etc., are small.

The effect of P odd admixtures in the wave functions of deuteron and the two nucleon continuum to the asymmetry is now shown to be quite small in the kinematic region of these quasi-elastic scattering experiments [24, 25]. However, it was realised that corrections due to the related processes of elastic electron deuteron scattering and coherent pion production in the kinematic region of this experiment could be important. It was also emphasized that quasi-elastic electron deuteron scattering at backward angles is predominantly sensitive to  $G_A^S$  and thus can be used to determine  $G_A^S$ . With a better knowledge of  $G_A^S$  from deuterium experiments, the extraction of  $G_E^S$  and  $G_M^S$  from hydrogen measurements will be facilitated. In view of this, the backward angle measurements of the parity violating asymmetry are being done in various scattering experiments with polarized electrons on deuterium targets [26]-[27].

In theory, many calculations have been done for the parity violating helicity dependent electron asymmetry in quasi-elastic scattering of polarized electrons off deuterons [24, 25] and [28]-[34]. But there exist very few calculations

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for the case of elastic electron deuteron scattering [34]-[37]. Recently, we have presented a general formalism for calculating all parity violating observables in elastic scattering of polarized and unpolarized electrons from polarized and/or unpolarized deuteron targets due to the interference of the weak and electromagnetic amplitudes [38]. The contribution of the strangeness component of the nucleon to the electron asymmetry in the scattering of polarized electrons from deuteron along with other contributions from the P odd parity admixture in the deuteron due to a parity violating nucleon-nucleon potential, anapole moments, exchange currents, and radiative corrections to the weak coupling constants are important theoretical ingredients, which should be studied in order to analyse the present experiments being done to observe the parity violating asymmetries in scattering experiments with polarized and unpolarized electrons.

In this paper we have studied the parity violating observables in elastic scattering of polarized and unpolarized electrons from unpolarized deuteron targets at electron energies relevant for SAMPLE [15, 16], PVA4 [26] and G0 [27] experiments. In particular, we have calculated the nonvanishing parity violating parts of the deuteron vector recoil polarisation for elastic scattering of unpolarized electrons from unpolarized deuterons. In the case of elastic scattering of polarized electrons from unpolarized deuterons, we have calculated the parity violating helicity dependent electron asymmetry as well as the parity violating nonvanishing components of the recoil deuteron polarisation. The separate contributions due to the  $\gamma$ - $Z$  interference and the odd parity admixture in the deuteron arising from the parity violating nucleon-nucleon potential have been calculated. The additional contribution due to the strangeness components of the nucleon, including anapole moments and radiative corrections to the weak coupling constants have been evaluated. The effect of meson exchange currents in the vector current contribution has been included through the use of Siegert's theorem for calculating the parity violating electric amplitude due to the vector current which is non-vanishing because of the odd parity admixture in the deuteron wave function.

In section II, we describe briefly the basic formalism for calculating all the parity violating observables in elastic scattering of polarized electrons from unpolarized deuterons. In section III and IV, we describe some technical details regarding the odd parity wave function components of the deuteron and the electroweak currents needed for calculating various contributions to parity violating observables. In section V, we present the numerical results, and a summary of our work with conclusions given in section VI.

## II. FORMALISM

Here we will give a brief review of the salient features of the formal expressions for cross section and polarization observables starting with a short outline of the classification of the various hadronic currents as defined in [38] whose multipoles enter in the expressions for the observables given below in subsection II C.

### A. Classification of Currents

The starting point of the classification is the general expression for the invariant scattering matrix elements containing both contributions from virtual  $\gamma$  and  $Z$  exchange

$$\mathcal{M}_{fi} = \frac{e^2}{q_\mu^2} \left( j_{fi}^{\gamma, \mu} J_{fi, \mu}^\gamma + \frac{1}{\sin^2(2\theta_W)} \frac{q_\mu^2}{M_Z^2 - q_\mu^2} j_{fi}^{Z, \mu} J_{fi, \mu}^Z \right), \quad (1)$$

denoting the electromagnetic and neutral currents of lepton and hadron by  $j_\mu^{\gamma/Z}$  and  $J_\mu^{\gamma/Z}$ , respectively, and with  $\theta_W$  the Weinberg angle. The lepton electromagnetic and weak neutral currents are separated into vector ( $j^v$ ) and axial vector ( $j$ ) contributions according to

$$j^{\gamma, \mu} = \bar{u}(k_2) \gamma^\mu u(k_1) = j^{v, \mu}, \quad (2)$$

$$j^{Z, \mu} = \bar{u}(k_2) (g_v \gamma^\mu + g_a \gamma^\mu \gamma_5) u(k_1) = j^{v, \mu} + j^{a, \mu}. \quad (3)$$

With respect to the hadron current, it is useful for the evaluation to distinguish between the contribution which couples to the lepton vector current and the one coupling to the lepton axial vector current by writing the matrix element in the form

$$\mathcal{M}_{fi} = \frac{e^2}{q_\mu^2} \left( j^{v, \mu} J_{fi, \mu}(\mathcal{V}) + j^{a, \mu} J_{fi, \mu}(\mathcal{A}) \right) \quad (4)$$

where we have introduced

$$J_{fi,\mu}(\mathcal{V}) = J_{fi,\mu}^\gamma + J_{fi,\mu}^{Z^V}, \quad (5)$$

$$J_{fi,\mu}(\mathcal{A}) = J_{fi,\mu}^{Z^A}, \quad (6)$$

with

$$J_{fi,\mu}^{Z^{\mathcal{V}/\mathcal{A}}} = \tilde{G}_{v/a} J_{fi,\mu}^Z, \quad (7)$$

where

$$\tilde{G}_v = \frac{g_v}{\sin^2(2\theta_W)} \frac{q_\mu^2}{M_Z^2 - q_\mu^2} = \frac{4 \sin^2 \theta_W - 1}{2 \sin^2(2\theta_W)} \frac{q_\mu^2}{M_Z^2 - q_\mu^2}, \quad (8)$$

$$\tilde{G}_a = \frac{g_a}{\sin^2(2\theta_W)} \frac{q_\mu^2}{M_Z^2 - q_\mu^2} = \frac{1}{2 \sin^2(2\theta_W)} \frac{q_\mu^2}{M_Z^2 - q_\mu^2}. \quad (9)$$

Here, the arguments  $\mathcal{V}$  and  $\mathcal{A}$  merely indicate to which type of lepton current the hadronic current couples. Both current types contain vector as well as axial contributions. We have further classified in [38] both contributions by their vector and axial parts. The electromagnetic current consists of only a vector piece while the neutral current contains both vector and axial ones, i.e. in an obvious notation

$$J_{fi,\mu}^Z = J_{fi,\mu}^{Z_v} + J_{fi,\mu}^{Z_a}. \quad (10)$$

Thus we have

$$\begin{aligned} J_{fi,\mu}(\mathcal{V}/\mathcal{A}) &= J_{fi,\mu}^\gamma + \tilde{G}_{v/a}(J_{fi,\mu}^{Z_v} + J_{fi,\mu}^{Z_a}) \\ &= J_{fi,\mu}^\gamma + J_{fi,\mu}^{Z^{\mathcal{V}/\mathcal{A}}} + J_{fi,\mu}^{Z^{\mathcal{V}/\mathcal{A}}}. \end{aligned} \quad (11)$$

Altogether, we have five types of hadronic currents, namely three of vector type  $J_{fi,\mu}^\gamma$ ,  $J_{fi,\mu}^{Z^V}$ , and  $J_{fi,\mu}^{Z^A}$ , and two of axial vector type  $J_{fi,\mu}^{Z^V}$  and  $J_{fi,\mu}^{Z^A}$ .

## B. Definition of Multipoles

The multipole moments appearing in the expressions of the various observables below are defined by the multipole expansion of the  $t$ -matrix contributions to the elastic scattering process arising from the hadronic vector and axial vector current contributions according to the classification in the previous subsection for which we will use a common label “ $c$ ”

$$\begin{aligned} t_{m'\lambda m}^c &= \frac{\sqrt{E'_d E_d}}{M_d} \langle m' | J_\lambda(c) | m \rangle \\ &= \frac{\sqrt{E'_d E_d}}{M_d} (-)^{\lambda} a_\lambda \sum_L i^L \hat{L} \langle 1m' | \mathcal{O}_{L\lambda}(c) | 1m \rangle \\ &= (-)^{1-m'+\lambda} a_\lambda \sum_L i^L \hat{L} \begin{pmatrix} 1 & L & 1 \\ -m' & \lambda & m \end{pmatrix} \mathcal{O}_L^\lambda(c), \end{aligned} \quad (12)$$

where  $a_\lambda = \sqrt{2\pi(1+\delta_{\lambda 0})}$  and the general multipole operator is defined by

$$\mathcal{O}_{LM}^\lambda(c) = \delta_{\lambda 0} \mathcal{C}_{LM}(c) + \delta_{|\lambda|1} (\mathcal{E}_{LM}(c) + \lambda \mathcal{M}_{LM}(c)), \quad (13)$$

with  $\mathcal{C}_{LM}$ ,  $\mathcal{E}_{LM}$ , and  $\mathcal{M}_{LM}$  denoting charge, electric and magnetic multipoles, respectively. The corresponding reduced matrix elements between deuteron states are given by

$$\begin{aligned} \mathcal{O}_L^\lambda(c) &= \frac{\sqrt{E'_d E_d}}{M_d} \langle 1 | \mathcal{O}_L^\lambda(c) | 1 \rangle \\ &= \delta_{\lambda 0} C_L(c) + \delta_{|\lambda|1} (E_L(c) + \lambda M_L(c)). \end{aligned} \quad (14)$$

### C. Observables

We will start from the general expression for a polarization observable of elastic electron deuteron scattering

$$e(k_1) + d(d) \rightarrow e(k_2) + d(d') \quad (15)$$

including longitudinal electron polarization of degree  $h$  but for an unpolarized deuteron target as derived in [38] and given there in Eq. (123)

$$\mathcal{O}_X \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \sigma_{\text{Mott}} S_0 \left[ A_d^0(X) + h A_{ed}^0(X) \right], \quad (16)$$

with

$$\sigma_{\text{Mott}} = \frac{\alpha^2 \cos^2 \frac{\theta_e^{\text{lab}}}{2}}{4 \sin^4 \frac{\theta_e^{\text{lab}}}{2}} \frac{k_2^{\text{lab}}}{(k_1^{\text{lab}})^3} \quad (17)$$

for the Mott cross section.

The four-momenta of incoming and scattered electrons are denoted by  $k_1$  and  $k_2$ , respectively, and the corresponding deuteron four-momenta by  $d = (E_d, \vec{d})$  and  $d' = (E'_d, \vec{d}')$ , respectively. Furthermore,  $q_\mu^2 = q_0^2 - \vec{q}^2$  denotes the squared four-momentum transfer with  $q = k_1 - k_2$ . The coordinate system chosen is such that the  $z$ -axis is taken along the momentum transfer  $\vec{q}$ , the  $y$ -axis along  $\vec{k}_1 \times \vec{k}_2$ , i.e. perpendicular to the scattering plane, and the  $x$ -axis as to form a right-handed system.

For an observable  $\mathcal{O}_X$  we had introduced in [38] the short-hand notation  $X = (IM\pm)$  with  $I = 0, 1, 2$  and  $I \geq M \geq 0$ . In detail, for  $I = 0$  one has only  $(00+)$  which refers to the differential cross section,  $(1M\pm)$  to the deuteron vector and  $(2M\pm)$  the tensor recoil polarization components in the spherical basis. The cartesian components are given for the vector polarization by

$$P_{x/y} = \pm \frac{1}{\sqrt{3}} \mathcal{O}_{11\pm}, \quad P_z = \sqrt{\frac{2}{3}} \mathcal{O}_{10+}, \quad (18)$$

and for tensor polarization

$$\begin{aligned} P_{xx/yy} &= \pm \frac{1}{2\sqrt{3}} \mathcal{O}_{22+} - \frac{1}{2} P_{zz}, & P_{zz} &= \frac{\sqrt{2}}{3} \mathcal{O}_{20+}, \\ P_{xy} &= -\frac{1}{2\sqrt{3}} \mathcal{O}_{22-}, & P_{zx/zy} &= \mp \frac{1}{2\sqrt{3}} \mathcal{O}_{21\pm}, \end{aligned} \quad (19)$$

where the cartesian components are defined by the deuteron density matrix in the form

$$\rho^d = \frac{1}{3} \left( 1 + \vec{P} \cdot \vec{S} + \sum_{kl} P_{kl} S_{kl}^{[2]} \right). \quad (20)$$

Here  $\vec{S}$  denotes the deuteron spin operator and  $S^{[2]}$  the corresponding irreducible tensor of rank 2

$$S_{kl}^{[2]} = \frac{1}{2} (S_k S_l + S_l S_k) - \frac{1}{3} \delta_{kl}. \quad (21)$$

We now will list the various asymmetries of (16).

(i) Differential cross section:

$$\frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \sigma_{\text{Mott}} S_0 \left[ 1 + A_d^{00+}(00+)_{pv} + h A_{ed}^{00+}(00+)_{pv} \right], \quad (22)$$

denoting by the subscript “ $pv$ ” a parity violating contribution. The unpolarized parity conserving differential cross section is given by  $\sigma_{\text{Mott}} S_0$  with

$$S_0 = \frac{4\pi}{3} \left( ((C_0^\gamma)^2 + (C_2^\gamma)^2) v_L + (M_1^\gamma)^2 v_T \right), \quad (23)$$

and the parity violating asymmetries by

$$S_0 A_d^{00+}(00+)_{pv} = \frac{8\pi}{3} E_1^{Z_a^A} M_1^\gamma v'_T, \quad (24)$$

$$S_0 A_{ed}^{00+}(00+)_{pv} = \frac{8\pi}{3} [(E_1^\gamma + E_1^{Z_a^V}) M_1^\gamma v'_T + (C_0^\gamma C_0^{Z_v^A} + C_2^\gamma C_2^{Z_v^A}) v_L + M_1^\gamma M_1^{Z_v^A} v_T]. \quad (25)$$

The kinematic functions  $v_\alpha^{(\prime)}$  ( $\alpha \in |L, T, LT, TT\rangle$ ) reflecting the virtual photon density matrix are in detail

$$\begin{aligned} v_L &= \beta^2 \xi^2, & v_T &= \frac{1}{2}(2\zeta + \xi), \\ v_{LT} &= 2\beta\xi, & v_{TT} &= -\frac{1}{2}\xi, \\ v'_{LT} &= \beta\sqrt{2\zeta}\xi, & v'_T &= \sqrt{\zeta(\zeta + \xi)}, \end{aligned} \quad (26)$$

with

$$\beta = \frac{|\vec{q}^{\text{lab}}|}{|\vec{q}^c|}, \quad \xi = -\frac{q_\nu^2}{|\vec{q}^{\text{lab}}|^2}, \quad \zeta = \tan^2 \frac{\theta_e^{\text{lab}}}{2}, \quad (27)$$

where  $\beta$  expresses the boost from the lab system to the frame in which the multipoles are evaluated and  $\vec{q}^c$  denotes the three-momentum transfer in this frame.

Collecting all terms one obtains

$$\begin{aligned} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} &= \frac{4\pi}{3} \sigma_{\text{Mott}} \left[ ((C_0^\gamma)^2 + (C_2^\gamma)^2) v_L + (M_1^\gamma)^2 v_T + 2E_1^{Z_a^A} M_1^\gamma v'_T \right. \\ &\quad \left. + 2h \left( (E_1^\gamma + E_1^{Z_a^V}) M_1^\gamma v'_T + (C_0^\gamma C_0^{Z_v^A} + C_2^\gamma C_2^{Z_v^A}) v_L + M_1^\gamma M_1^{Z_v^A} v_T \right) \right], \end{aligned} \quad (28)$$

where  $C_L^c$ ,  $E_L^c$ , and  $M_L^c$  are the charge and transverse electric and magnetic multipoles, respectively, defined in subsection II B for various hadronic currents  $c$  ( $\gamma$ ,  $Z_v$ , or  $Z_a$ ) and which are given explicitly below in section IV A.

In the Standard Model without radiative corrections and in the absence of the strangeness contribution to the neutral current, all multipole matrix elements are proportional to the conventional electromagnetic ones. According to the relations given in [38], we get for isoscalar transitions

$$E_1^{Z_a^A} = \tilde{G}_a F_{E1}^A, \quad E_1^{Z_a^V} = \tilde{G}_v F_{E1}^A, \quad \text{where } F_{E1}^A = \frac{\sqrt{E_d'E_d}}{M_d}, \langle 1 | E_1 (J^{Z_a}) | 1 \rangle, \quad (29)$$

$$M_1^{Z_v^A} = g_v^d \tilde{G}_a M_1^\gamma, \quad C_L^{Z_v^A} = g_v^d \tilde{G}_a C_L^\gamma, \quad \text{where } g_v^d = -2 \sin^2 \theta_W. \quad (30)$$

Furthermore,  $C_L^\gamma$  ( $L = 0, 2$ ) and  $M_1^\gamma$  denote the usual electromagnetic deuteron charge and magnetic multipoles,  $F_{E1}^A$  its electric multipole due to axial current and  $E_1^\gamma$  the electric multipole due to a parity violating admixture in the wave function. Thus, in the Standard Model the cross section becomes

$$\begin{aligned} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} &= \frac{4\pi}{3} \sigma_{\text{Mott}} \left[ (1 - h4 \sin^2 \theta_W \tilde{G}_a) \left( (C_0^\gamma)^2 + (C_2^\gamma)^2 \right) v_L + (M_1^\gamma)^2 v_T \right. \\ &\quad \left. + 2 \left( h E_1^\gamma + (\tilde{G}_a + h \tilde{G}_v) F_{E1}^A \right) M_1^\gamma v'_T \right]. \end{aligned} \quad (31)$$

(ii) Vector recoil polarization:

$$P_x \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = -\frac{1}{\sqrt{3}} \sigma_{\text{Mott}} S_0 \left( A_d^{00+}(11+)_{pv} + h A_{ed}^{00+}(11+)_{pc+pv} \right), \quad (32)$$

$$P_y \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = 0, \quad (33)$$

$$P_z \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \sqrt{\frac{2}{3}} \sigma_{\text{Mott}} S_0 \left( A_d^{00+}(10+)_{pv} + h A_{ed}^{00+}(10+)_{pc+pv} \right), \quad (34)$$

where the asymmetry parameters  $A_d^{00+}(11+)_{pv}$ ,  $A_d^{00+}(10+)_{pv}$ ,  $A_{ed}^{00+}(11+)_{pc+pv}$ , and  $A_{ed}^{00+}(10+)_{pc+pv}$  are given by

$$\begin{aligned} S_0 A_d^{00+}(11+)_{pv} &= \frac{4\pi}{3} \left( E_1^\gamma + E_1^{Z_a^V} \right) \left( 2\sqrt{2}C_0^\gamma + C_2^\gamma \right) v_{LT} \\ &\quad + \frac{2\sqrt{2}}{3} \left[ \left( 4C_0^{Z_v^A} + \sqrt{2}C_2^{Z_v^A} \right) M_1^\gamma + \left( 4C_0^\gamma + \sqrt{2}C_2^\gamma \right) M_1^{Z_v^A} \right] v'_{LT}, \end{aligned} \quad (35)$$

$$S_0 A_{ed}^{00+}(11+)_{pc} = \frac{2\sqrt{2}\pi}{3} \left( 4C_0^\gamma + \sqrt{2}C_2^\gamma \right) M_1^\gamma v'_{LT}, \quad (36)$$

$$S_0 A_{ed}^{00+}(11+)_{pv} = \frac{4\pi}{3} \left( 2\sqrt{2}C_0^\gamma + C_2^\gamma \right) E_1^{Z_a^A} v_{LT}, \quad (37)$$

$$S_0 A_d^{00+}(10+)_{pv} = 2\sqrt{\frac{2}{3}}\pi \left[ \left( E_1^\gamma + E_1^{Z_a^V} \right) M_1^\gamma v_T + M_1^\gamma M_1^{Z_v^A} v'_T \right], \quad (38)$$

$$S_0 A_{ed}^{00+}(10+)_{pc} = \sqrt{\frac{2}{3}}\pi (M_1^\gamma)^2 v'_T \quad (39)$$

$$S_0 A_{ed}^{00+}(10+)_{pv} = 2\sqrt{\frac{2}{3}}\pi E_1^{Z_a^A} M_1^\gamma v_T. \quad (40)$$

Introducing for convenience for any of the current contributions

$$C^c = 4C_0^c + \sqrt{2}C_2^c, \quad (41)$$

the evaluation of the asymmetries yields

$$\begin{aligned} P_x \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} &= -\frac{4\pi}{3\sqrt{6}} \sigma_{\text{Mott}} \left[ (E_1^\gamma + E_1^{Z_a^V}) C^\gamma v_{LT} + (C^{Z_v^A} M_1^\gamma + C^\gamma M_1^{Z_v^A}) v'_{LT} \right. \\ &\quad \left. + h(M_1^\gamma v'_{LT} + E_1^{Z_a^A} v_{LT}) C^\gamma \right], \end{aligned} \quad (42)$$

$$P_z \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \frac{2\pi}{3} \sigma_{\text{Mott}} \left[ 2(E_1^\gamma + E_1^{Z_a^V}) v_T + 2M_1^{Z_v^A} v'_T + h(M_1^\gamma v'_T + 2E_1^{Z_a^V} v_T) \right] M_1^\gamma. \quad (43)$$

Using the lepton coupling from Eqs. (29) and (30) one has in the Standard Model

$$P_x \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = -\frac{4\pi}{3\sqrt{6}} \sigma_{\text{Mott}} \left[ \left( E_1^\gamma + (\tilde{G}_v + h\tilde{G}_a) F_{E1}^A \right) v_{LT} + (h - 2\sin^2 \theta_W \tilde{G}_a) M_1^\gamma v'_{LT} \right] C^\gamma, \quad (44)$$

$$P_z \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \frac{2\pi}{3} \sigma_{\text{Mott}} \left[ 2 \left( E_1^\gamma + (\tilde{G}_v + h\tilde{G}_a) F_{E1}^A \right) v_T + (h - 4\sin^2 \theta_W \tilde{G}_a) M_1^\gamma v'_T \right] M_1^\gamma. \quad (45)$$

We would like to point out, that the nonvanishing contribution to the vector recoil polarization for unpolarized electrons is solely due to parity violating contributions (see Eqs. (32)-(34)).

(ii) Tensor recoil polarization:

The nonvanishing components are

$$P_{zz} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \frac{\sqrt{2}}{3} \sigma_{\text{Mott}} S_0 \left( A_d^{00+}(20+)_{pc+pv} + h A_{ed}^{00+}(20+)_{pv} \right), \quad (46)$$

$$P_{xx/yy} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \pm \frac{1}{2\sqrt{3}} \sigma_{\text{Mott}} S_0 \left( A_d^{00+}(22+)_{pc+pv} + h A_{ed}^{00+}(22+)_{pv} \right) - \frac{1}{2} P_{zz} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}}, \quad (47)$$

$$P_{zx} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = -\frac{1}{2\sqrt{3}} \sigma_{\text{Mott}} S_0 \left( A_d^{00+}(21+)_{pc+pv} + h A_{ed}^{00+}(21+)_{pv} \right), \quad (48)$$

where the parameters  $A_d^{00+}(20+)_{pc+pv}$ ,  $A_d^{00+}(21+)_{pc+pv}$ ,  $A_{ed}^{00+}(20+)_{pv}$ ,  $A_{ed}^{00+}(21+)_{pv}$  and  $A_{ed}^{00+}(22+)_{pv}$  are given

by

$$S_0 A_d^{00+}(20+)_{pc} = -\frac{2\pi}{3} \left( 4C_0^\gamma + \sqrt{2}C_2^\gamma \right) C_2^\gamma v_L - \frac{\sqrt{2}\pi}{3} (M_1^\gamma)^2 v_T, \quad (49)$$

$$S_0 A_d^{00+}(20+)_{pv} = -\frac{2\sqrt{2}}{3} \pi E_1^{Z_a^A} M_1^\gamma v'_T, \quad (50)$$

$$S_0 A_d^{00+}(21+)_{pc} = 4\pi C_2^\gamma M_1^\gamma v_{LT}, \quad (51)$$

$$S_0 A_d^{00+}(21+)_{pv} = 4\pi C_2^\gamma E_1^{Z_a^A} v'_{LT}, \quad (52)$$

$$\begin{aligned} S_0 A_{ed}^{00+}(20+)_{pv} &= -\frac{4\pi}{3} \left( 2C_0^{Z_v^A} C_2^\gamma + 2C_0^\gamma C_2^{Z_v^A} + \sqrt{2}C_2^\gamma C_2^{Z_v^A} \right) v_L \\ &\quad - \frac{2\sqrt{2}}{3} \pi \left[ (M_1^\gamma M_1^{Z_v^A}) v_T + (E_1^\gamma + E_1^{Z_a^V}) M_1^\gamma v'_T \right], \end{aligned} \quad (53)$$

$$S_0 A_{ed}^{00+}(21+)_{pv} = 4\pi \left[ (C_2^{Z_v^A} M_1^\gamma + C_2^\gamma M_1^{Z_v^A}) v_{LT} + (E_1^\gamma + E_1^{Z_a^V}) C_2^\gamma v'_{LT} \right], \quad (54)$$

$$S_0 A_{ed}^{00+}(22+)_{pv} = \frac{4\pi}{\sqrt{3}} M_1^\gamma M_1^{Z_v^A} v_{TT}. \quad (55)$$

Using these expressions, one gets

$$\begin{aligned} P_{zz} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} &= -\frac{2\pi}{9} \sigma_{\text{Mott}} \left[ \sqrt{2}C^\gamma C_2^\gamma v_L + (M_1^\gamma)^2 v_T + 2E_1^{Z_a^A} M_1^\gamma v'_T \right. \\ &\quad \left. - 2h \left( \sqrt{2}(2C_0^\gamma C_2^{Z_v^A} + 2C_0^{Z_v^A} C_2^\gamma + \sqrt{2}C_2^\gamma C_2^{Z_v^A}) v_L + (M_1^{Z_v^A} v_T + (E_1^\gamma + E_1^{Z_a^V}) v'_T) M_1^\gamma \right) \right], \end{aligned} \quad (56)$$

$$P_{xx/yy} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \pm \frac{\pi}{3} \sigma_{\text{Mott}} \left[ M_1^\gamma + h2M_1^{Z_v^A} \right] M_1^\gamma v_{TT} - \frac{1}{2} P_{zz} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}}, \quad (57)$$

$$\begin{aligned} P_{zx} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} &= -\frac{2\pi}{\sqrt{3}} \sigma_{\text{Mott}} \left[ C_2^\gamma M_1^\gamma v_{LT} + E_1^{Z_a^A} C_2^\gamma v'_{LT} \right. \\ &\quad \left. + h \left( (C_2^{Z_a^V} M_1^\gamma + C_2^\gamma M_1^{Z_v^A}) v_{LT} + (E_1^\gamma + E_1^{Z_a^V}) C_2^\gamma v'_{LT} \right) \right]. \end{aligned} \quad (58)$$

Again, these expressions reduce in the Standard Model to

$$\begin{aligned} P_{zz} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} &= -\frac{2\pi}{9} \sigma_{\text{Mott}} \left[ (1 + h4\sin^2\theta_W \tilde{G}_a) \left( \sqrt{2}C^\gamma C_2^\gamma v_L + (M_1^\gamma)^2 v_T \right) \right. \\ &\quad \left. + 2 \left( (\tilde{G}_a - h\tilde{G}_v) F_{E1}^A - hE_1^\gamma \right) v'_T \right], \end{aligned} \quad (59)$$

$$P_{xx/yy} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = \pm \frac{\pi}{3} \sigma_{\text{Mott}} (1 - h4\sin^2\theta_W \tilde{G}_a) (M_1^\gamma)^2 v_{TT} - \frac{1}{2} P_{zz} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}}, \quad (60)$$

$$P_{zx} \frac{d\sigma^{\gamma+Z}}{d\Omega_{k_2}^{\text{lab}}} = -\frac{2\pi}{\sqrt{3}} \sigma_{\text{Mott}} \left[ (1 - h4\sin^2\theta_W \tilde{G}_a) M_1^\gamma v_{LT} + (E_1^\gamma + (\tilde{G}_a + h\tilde{G}_v) F_{E1}^A) v'_{LT} \right] C_2^\gamma. \quad (61)$$

### III. DEUTERON WAVE FUNCTIONS

#### A. Deuteron Parity Conserving Components

The isospin singlet parity conserving component of the deuteron wave function  $\psi_d^{pc}(\vec{r})$  is written as

$$\psi_d^{pc}(\vec{r}) = \sum_{L=0,2} i^L \frac{u_L(r)}{r} Y_{1L,1}^{m_d}(\hat{r}), \quad (62)$$

where

$$Y_{1L,1}^{m_d}(\hat{r}) = \sum_{m_L, m_S} \langle L m_L | 1 m_S \rangle Y_{LML}(\hat{r}) \chi_{1m_S}, \quad (63)$$

with  $\chi_{1m_S}$  denoting the deuteron spin wave function, and the radial parts  $u_L(r)$  are taken from the parameterization of Machleidt et al. [39]. for the Bonn OBEPQ model, given by

$$u_0(r) = \sum_{i=1}^{n_0} C_i e^{-m_i r}, \quad u_2(r) = \sum_{i=1}^{n_2} D_i e^{-m_i r} \left( 1 + \frac{3}{m_i r} + \frac{3}{(m_i r)^2} \right), \quad (64)$$

with the normalization

$$\sum_{L=0,2} \int_0^\infty u_L^2(r) dr = 1. \quad (65)$$

The constants  $C_i$  and  $m_i$  in Eq. (64) are listed in [39].

## B. Deuteron Parity Violating Components

The calculation of the small parity violating component in the deuteron due to the weak parity violating hadronic potential  $V^{pnc}$  is done in first order perturbation theory as described in [31] using the potential of [40]

$$\begin{aligned} V^{pnc}(\vec{r}, \vec{p}) = & i \frac{f_\pi g_{\pi NN}}{2\sqrt{2}M} (\vec{\tau}_1 \times \vec{\tau}_2)_z (\vec{\sigma}_1 + \vec{\sigma}_2) \cdot [\vec{p}, f_\pi(r)] \\ & - \frac{g_\rho}{M} \left( h_\rho^0 \vec{\tau}_1 \cdot \vec{\tau}_2 + \frac{h_\rho^1}{2} (\vec{\tau}_1 + \vec{\tau}_2)_z + \frac{h_\rho^2}{2\sqrt{6}} (3\tau_{1,z}\tau_{2,z} - \vec{\tau}_1 \cdot \vec{\tau}_2) \right) \\ & \times \left( (\vec{\sigma}_1 - \vec{\sigma}_2) \cdot \{ \vec{p}, f_\rho(r) \} + i(1 + \chi_v) (\vec{\sigma}_1 \times \vec{\sigma}_2) \cdot [\vec{p}, f_\rho(r)] \right) \\ & - \frac{g_\omega}{M} \left( h_\omega^0 + \frac{h_\omega^1}{2} (\vec{\tau}_1 + \vec{\tau}_2)_z \right) \\ & \times \left( (\vec{\sigma}_1 - \vec{\sigma}_2) \cdot \{ \vec{p}, f_\omega(r) \} + i(1 + \chi_s) (\vec{\sigma}_1 \times \vec{\sigma}_2) \cdot [\vec{p}, f_\omega(r)] \right) \\ & - \frac{1}{2M} (\vec{\tau}_1 - \vec{\tau}_2)_z (\vec{\sigma}_1 + \vec{\sigma}_2) \cdot \{ \vec{p}, g_\omega h_\omega^1 f_\omega(r) - g_\rho h_\rho^1 f_\rho(r) \} \\ & - i \frac{g_\rho h_\rho'^1}{2M} (\vec{\tau}_1 \times \vec{\tau}_2)_z (\vec{\sigma}_1 + \vec{\sigma}_2) \cdot [\vec{p}, f_\rho(r)], \end{aligned} \quad (66)$$

with the usual Yukawa function

$$f_\xi(r) = \frac{e^{-m_\xi r}}{4\pi r}, \quad \text{for } \xi = \pi, \rho, \omega. \quad (67)$$

$M$  denotes the nucleon mass and  $\vec{p} = \frac{1}{2}(\vec{p}_1 - \vec{p}_2)$ . Values for the weak coupling constants for various models of parity violating hadronic potentials are listed in Table I.

TABLE I: Weak coupling constants for various models of parity violating  $NN$  interaction  $V^{pnc}$  in units of  $g_\pi = 3.8 \times 10^{-8}$ .

Coupling	$f_\pi$	$h_\rho^0$	$h_\rho^1$	$h_\rho^2$	$h_\omega^0$	$h_\omega^1$	$h_\rho'^1$
DDH [40]	12	-30	-0.5	-25	-5	-3	0
DZ [41]	+3	-22	+1	-18	-10	-6	0
FCDH [42]	+7	-10	-1	-18	-13	-6	0

The small parity violating component in the deuteron wave function  $\psi_d^{pnc}$  is given in first order perturbation theory by

$$| \psi_d^{pnc} \rangle = - \frac{1}{H_{pc}^{(0)} - E} \left( H_{pnc}^{(1)} - (E - E_B) \right) | \psi_d^{pc} \rangle, \quad (68)$$

where  $| \psi_d^{pc} \rangle$  is the unperturbed parity conserving deuteron wave function,  $H_{pc}^{(0)}$  and  $H_{pnc}^{(1)}$  are the parity conserving and violating strong Hamiltonians, respectively,  $E_B$  is the unperturbed deuteron binding energy ( $= -2.2246$  MeV) and

$E$  the eigenvalue of the perturbed wave function. In the present calculation, the propagator was approximated by the free Greens function. Applying the potential of Eq. (66) to the unperturbed wave function yields for the parity violating admixture

$$\psi_d^{pnc}(\mathbf{p}) = \frac{i}{p} \left( \tilde{u}_{11}(p) \langle \hat{p}|10; (11)1m_d \rangle + \tilde{u}_{10}(p) \langle \hat{p}|00; (10)1m_d \rangle \right), \quad (69)$$

where  $\tilde{u}_{1S}(p)$  denotes the radial part and  $\langle \hat{p}|Tm_T; (1S)JM \rangle$  the isospin, orbital and spin angular momentum parts of the parity violating p-wave ( $L=1$ ) components of the deuteron. The two contributions  $u_{11}(p)$  and  $u_{10}(p)$  correspond to isovector  $^3P_1$  and isoscalar  $^1P_1$  states, respectively. In detail one finds

$$\tilde{u}_{1S}(p) = \frac{u_{1S}(p)}{p} = \frac{1}{E_B - \frac{\vec{p}^2}{M}} \int dr j_1(pr) f_{1S}(r), \quad E_B < 0, \quad (70)$$

where for the triplet state  $\tilde{u}_{11}(p)$  the function  $f_{11}(r)$  is given by

$$\begin{aligned} f_{11}(r) = & -\frac{1}{\pi M \sqrt{3\pi}} \sum_{L=0,2} (\sqrt{2})^{-L/2} \left\{ \frac{f_\pi g_{\pi NN}}{\sqrt{2}} e^{-m_\pi r} \left( m_\pi + \frac{1}{r} \right) u_L(r) \right. \\ & - g_\rho h_\rho^1 e^{-m_\rho r} \left( m_\rho + \frac{1}{r} \right) u_L(r) - g_\omega h_\omega^1 e^{-m_\omega r} \left[ \left( m_\omega + (-)^{L/2} \frac{3}{r} \right) u_l(r) - 2u'_L(r) \right] \\ & \left. + g_\rho h_\rho^1 e^{-m_\rho r} \left[ \left( m_\rho + (-)^{L/2} \frac{3}{r} \right) u_l(r) - 2u'_L(r) \right] \right\}, \end{aligned} \quad (71)$$

and for the singlet state  $\tilde{u}_{10}(p)$

$$\begin{aligned} f_{10}(r) = & -\frac{1}{\pi M \sqrt{6\pi}} \sum_{L=0,2} (-\sqrt{2})^{-L/2} \left\{ -3g_\rho h_\rho^0 e^{-m_\rho r} \left( \left[ \left( m_\rho + \frac{1}{r} \right) \chi_v + \frac{(-2)^{1+L/2}}{r} \right] u_L(r) + 2u'_L(r) \right) \right. \\ & \left. + g_\omega h_\omega^0 e^{-m_\omega r} \left( \left[ \left( m_\omega + \frac{1}{r} \right) \chi_S + \frac{(-2)^{1+L/2}}{r} \right] u_L(r) + 2u'_L(r) \right) \right\}. \end{aligned} \quad (72)$$

In configuration space, the parity violating components  $u_{1S}(r)$  are obtained by taking the Fourier transform of the momentum space wave functions  $u_{1S}(p)$  given in Eq. (70). Explicitly they are written as [43]

$$\begin{aligned} \frac{u_{1S}(r)}{r} = & \sqrt{\frac{2}{\pi}} \int_0^\infty dp p^2 \frac{u_{1S}(p)}{p} j_1(pr) \\ = & -\sqrt{\frac{2}{\pi}} M \int_0^\infty dr' G_1(r, r') f_{1S}(r), \end{aligned} \quad (73)$$

where

$$G_1(r, r') = \int_0^\infty dp p^2 \frac{j_1(pr) j_1(pr')}{p^2 + \epsilon^2} = \frac{1}{rr'} H_1(r, r', \epsilon), \quad (74)$$

with

$$\epsilon^2 = -E_B M \quad \text{and} \quad H_1(r, r', t) = \epsilon k_1(\epsilon r_>) i_1(\epsilon r_<), \quad (75)$$

where  $k_1$  and  $i_1$  are the modified Bessel and Hankel functions. For a given parity conserving wave function of the deuteron, the p-wave components are numerically quite sensitive to the parameters of the weak nucleon-nucleon potential, especially to  $f_\pi$ . These p-wave components are shown in Fig. 1 for various potential models whose parameters values are given in Table I.

#### IV. ELECTROWEAK CURRENTS AND MULTipoles

##### A. Multipoles

We see from section II, that in elastic deuteron scattering with unpolarized deuteron targets, the various observables like differential cross section  $d\sigma/d\Omega$ , vector and tensor polarizations of the recoil deuteron, i.e.  $P_i$  ( $i = x, y, z$ ) and

$P_{ij}$  ( $i, j = x, y, z$ ), are given in terms of the quantities  $A_d^{00+}(IM)$  and  $A_{ed}^{00+}(IM)$ , with  $I = 0, 1, 2$  and  $0 < M < I$  for unpolarized and polarized electron scattering. These quantities are defined in terms of multipoles  $C_0^c, C_2^c, E_1^c$ , and  $M_1^c$  which are reduced matrix elements of the general multipole operators (see Eq. (14)) for a given current  $c$ , which may be the vector current ( $J^\gamma$ ) of the electromagnetic interaction due to one photon exchange or the weak neutral vector and axial currents ( $J^{Z_v}$  and  $J^{Z_a}$ ) due to  $Z$  exchange. It should be noted that in the presence of parity violation the nonvanishing multipole  $E_1^c$  arises from the weak axial vector current ( $J^{Z_a}$ ) and also from the electromagnetic current ( $J^\gamma$ ) when the parity violating component is included in the wave function of the deuteron. The general multipole operators are defined by [44].

$$C_{LM}^c = \int d\vec{x} \left[ j_L(qx) \vec{Y}_{LM}(\hat{x}) \right] J_0^c(\vec{x}), \quad (76)$$

$$E_{LM}^c = \frac{1}{q} \int d\vec{x} \vec{\nabla} \times \left[ j_L(qx) \vec{Y}_{LL1}^M(\hat{x}) \right] \cdot \vec{J}^c(\vec{x}), \quad (77)$$

$$M_{LM}^c = \int d\vec{x} \left[ j_L(qx) \vec{Y}_{LL1}^M(\hat{x}) \right] \cdot \vec{J}^c(\vec{x}), \quad (78)$$

where the currents  $J_\mu^c$  ( $c = \gamma, Z_v, Z_a$ ) are given by their single nucleon matrix elements in terms of the weak and electromagnetic form factors of the nucleons as follows

$$\langle p' | J_\mu^{\gamma,p(n)} | p \rangle = \bar{u}(p') \left[ F_1^{\gamma,p(n)}(Q^2) \gamma_\mu + i F_2^{\gamma,p(n)}(Q^2) \frac{\sigma_{\mu\nu}}{2M} q^\nu \right] u(p), \quad (79)$$

$$\langle p' | J_\mu^{Z_v,p(n)} | p \rangle = \langle p' | J_\mu^{Z_v,p(n)} + J_\mu^{Z_a,p(n)} | p \rangle, \quad (80)$$

with

$$\langle p' | J_\mu^{Z_v,p(n)} | p \rangle = \bar{u}(p') \left[ F_1^{Z,p(n)}(Q^2) \gamma_\mu + i F_2^{Z,p(n)}(Q^2) \sigma_{\mu\nu} \frac{q^\nu}{2M} \right] u(p), \quad (81)$$

$$\langle p' | J_\mu^{Z_a,p(n)} | p \rangle = \bar{u}(p') \left[ G_A^{Z,p(n)}(Q^2) \gamma_\mu \gamma_5 \right] u(p), \quad (82)$$

where  $F_1^{\gamma,Z}(Q^2)$  and  $F_2^{\gamma,Z}(Q^2)$  ( $Q^2 = -q_\mu^2$ ) are the Dirac form factors of the electromagnetic and weak neutral vector currents which are related to the Sachs form factors  $G_E^{\gamma,Z}(Q^2)$  and  $G_M^{\gamma,Z}(Q^2)$  as follows

$$G_E^{\gamma,Z}(Q^2) = F_1^{\gamma,Z}(Q^2) - \tau F_2^{\gamma,Z}(Q^2), \quad (83)$$

$$G_M^{\gamma,Z}(Q^2) = F_1^{\gamma,Z}(Q^2) + F_2^{\gamma,Z}(Q^2), \quad \tau = \frac{Q^2}{4M}. \quad (84)$$

The weak neutral current form factors  $G_{E,M}^Z(Q^2)$  are defined in the standard model [45] by

$$G_{E,M}^{Z,p(n)}(Q^2) = \frac{1}{2} (1 - 4 \sin^2 \theta_W) G_{E,M}^{p(n)}(Q^2) - \frac{1}{2} G_{E,M}^{n(p)}(Q^2), \quad (85)$$

where  $\theta_W$  is the weak mixing angle in the standard model. The axial vector form factors are given by

$$G_A^{Z,p(n)}(Q^2) = -\frac{\tau_3}{2} G_A^Z(Q^2), \quad \tau_3 = +1(-1) \text{ for } p(n). \quad (86)$$

In the presence of nonzero strangeness of the nucleon, these form factors are modified as

$$G_{E,M}^{Z,p(n)}(Q^2) \rightarrow G_{E,M}^{Z,p(n)}(Q^2) - \frac{1}{2} G_{E,M}^S(Q^2), \quad (87)$$

$$G_A^{Z,p(n)}(Q^2) \rightarrow G_A^{Z,p(n)}(Q^2) + \frac{1}{2} G_A^S(Q^2). \quad (88)$$

The matrix elements of the multipole operators between the initial and final deuteron state are calculated using the nonrelativistic limit of the current matrix elements for the vector and axial currents given by:

$$\left\{ J_0, \vec{J}_\gamma \right\} = \chi_{s_f}^\dagger \left\{ G_E^\gamma(Q^2), \frac{1}{2M} [G_E^\gamma(Q^2) (\vec{p}_f + \vec{p}_i) + i G_M^\gamma(Q^2) (\vec{\sigma} \times \vec{q})] \right\} \chi_{s_i}, \quad (89)$$

$$\left\{ J_0, \vec{J}_{Z_v} \right\} = \chi_{s_f}^\dagger \left\{ G_E^Z(Q^2), \frac{1}{2M} [G_E^Z(Q^2) (\vec{p}_f + \vec{p}_i) + i G_M^Z(Q^2) (\vec{\sigma} \times \vec{q})] \right\} \chi_{s_i}, \quad (90)$$

$$\left\{ J_0, \vec{J}_{Z_a} \right\} = \chi_{s_f}^\dagger \left\{ \frac{G_A^Z(Q^2)}{2M} \vec{\sigma} \cdot (\vec{p}_i + \vec{p}_f), -G_A^Z(Q^2) \vec{\sigma} \right\} \chi_{s_i}. \quad (91)$$

## B. Electroweak Form Factors

### 1. Electromagnetic and Weak Form Factors

The electromagnetic form factors  $G_{E,M}^{\gamma,p(n)}(Q^2)$  defined in Eqs. (83) and (84) are generally parameterized in dipole forms given by

$$G_E^p(Q^2) = \left[1 + \frac{Q^2}{M_V^2}\right]^{-2}, \quad G_E^n(Q^2) = \mu_n \xi_n \frac{Q^2}{M^2} G_E^p(Q^2), \quad (92)$$

$$G_M^p(Q^2) = (1 + \mu_p) G_E^p(Q^2), \quad G_M^n(Q^2) = \mu_n G_E^p(Q^2), \quad (93)$$

with

$$\xi_n = \left[1 - \lambda_n \left(\frac{Q^2}{4M^2}\right)\right]^{-1}, \quad \mu_p = 1.792847, \quad \mu_n = -1.913043, \quad \lambda_n = 5.6. \quad (94)$$

The axial vector form factor  $G_A^{Z,p(n)}(Q^2)$  is also parameterized in dipole form as

$$G_A^{Z,p(n)}(Q^2) = G_A^{Z,p(n)}(0) \left[1 + \frac{Q^2}{M_A^2}\right]^{-2} \text{ with } G_A^{Z,p(n)}(0) = -\frac{1}{2} \left(+\frac{1}{2}\right) 1.262. \quad (95)$$

The numerical value of the vector dipole mass  $M_V = 0.84$  GeV is taken from experimental data on electron proton scattering, and the axial dipole mass  $M_A = 1.026$  GeV from neutrino scattering from proton and deuteron [46].

### 2. Strangeness Form Factors

A number of theoretical models have been used to establish the magnitude of the strange quark form factor and its  $Q^2$  dependence. They use either a pole with simple vector dominance for the vector form factors or a meson cloud model which considers kaon loops and other higher strange resonance loop contributions to calculate the strangeness form factors. A review of many of these models is given in [47]. Recently, chiral perturbation theory [48] and lattice QCD [49, 50] have also been applied for their calculation. In the earlier theoretical work a dipole form has been used. However, in the present calculation we have used for the vector strangeness form factor the recently determined form factors from a global analysis of presently available data on parity violating electron scattering by Liu et al. [17], i.e. with  $Q^2$  in  $(\text{GeV}/c)^2$

$$G_E^s(Q^2) = G_E^s \frac{Q^2}{0.1}, \quad G_E^s = -0.014 \quad (96)$$

$$G_M^s(Q^2) = G_M^s + \mu'_s (Q^2 - 0.1), \quad G_M^s = 0.28, \quad \mu'_s = -0.1 (\text{GeV}/c)^{-2}. \quad (97)$$

For the axial vector form factor  $G_A^s(Q^2)$  we use a dipole form

$$G_A^s(Q^2) = g_A^s(0) \left[1 + \frac{Q^2}{M_A^2}\right]^{-2}, \quad (98)$$

with  $g_A^s(0) = \Delta s = -0.19$  and  $M_A = 1.026$  GeV, where  $\Delta s$  is the spin contribution of the strange quarks and antiquarks determined experimentally from deep inelastic scattering of electrons.

### 3. Radiative Corrections and Anapole Moments

Higher order radiative corrections to the weak neutral current couplings of the nucleon have been calculated in the standard model [19]-[23]. The corrections to the vector form factors are dominated by single quark transition but their effect is found to be small as they are multiplied by the factor  $(1 - 4 \sin^2 \theta_W)$ . In the presence of radiative corrections, the weak vector form factors are written as

$$G_{E,M}^{Z,p(n)}(Q^2) = \frac{1}{2} (1 - 4 \sin^2 \theta_W) (1 + R_V^p) G_{E,M}^p(Q^2) - \frac{1}{2} (1 + R_V^n) G_{E,M}^n(Q^2) - \frac{1}{2} G_{E,M}^s(Q^2), \quad (99)$$

where  $R_V^p$  and  $R_V^n$  are the radiative corrections for proton and neutron. In the case of the axial vector, the single quark and two quark transitions are both important and lead to appreciable corrections. The axial vector form factors in the presence of radiative correction are written as

$$G_A^Z(Q^2) = -\frac{1}{2}\tau_3(1+R_A^1)G_A - \frac{1}{2}R_A^0 + \frac{1}{2}G_A^s, \quad (100)$$

where  $R_A^1$  and  $R_A^0$  are the radiative corrections in the isovector and isoscalar channels. In addition to single quark transition the two quark transition induce the following axial anapole term in the matrix element of the electromagnetic current

$$\langle p' | J_\mu^\gamma | p \rangle = \frac{Q^2}{M^2} \bar{u}(p') \left[ \gamma_\mu - \frac{q q_\mu}{q^2} \right] \gamma_5 u(p) [a_S(Q^2) + a_V(Q^2)\tau_3], \quad (101)$$

which in leading order of the nonrelativistic limit is given as

$$\vec{J} = \frac{Q^2}{M^2} \chi_{s_f}^\dagger [\vec{\sigma} - \vec{\sigma} \cdot \hat{q} \hat{q}] \chi_{s_i} [a_S(Q^2) + a_V(Q^2)\tau_3] \quad (102)$$

with  $\hat{q}$  as unit vector along  $\vec{q}$ . It is equivalent to an axial coupling contributing to the axial form factor. The coefficients  $a_S(Q^2)$  and  $a_V(Q^2)$  are calculated using pion loop contributions in terms of the parity violating and parity conserving  $\pi NN$  couplings  $f_{\pi NN}$  and  $g_{\pi NN}$ , and are given as [21]

$$a_{S,V}(0) = \frac{f_{\pi NN} g_{\pi NN}}{4\sqrt{2}\pi^2} \alpha_{S,V}(0), \text{ with } \alpha_S(0) = 1.6, \alpha_V(0) = 0.4. \quad (103)$$

The contributions of these terms to the radiative corrections  $R_A^1$  and  $R_A^0$  are found to be small [21, 45]. The updated values of the radiative corrections, taken from ref. [7], are given in Table II.

TABLE II: Values of radiative corrections to weak neutral current couplings.

Correction	Isoscalar	Isovector
$R_V$	-0.0113	-0.017±0.002
$R_A$	0.06±0.14	-0.23±0.24

## V. RESULTS AND DISCUSSIONS

The differential cross sections  $d\sigma/d\Omega$ , the vector and tensor polarizations  $P_i$  ( $i = x, y, z$ ) and  $P_{ij}$  ( $i, j = x, y, z$ ), respectively, are described in terms of various asymmetry parameters  $A_d^{00+}(IM)$  and  $A_{ed}^{00+}(IM)$ . They get contributions from the parity conserving as well as from the parity and time reversal (T) violating pieces. In Table III, we give a classification of parity conserving and parity violating contributions to  $A_d^{00+}(IM)$ , and  $A_{ed}^{00+}(IM)$  for various values of  $I$  and  $M$ .

TABLE III: Schematic survey of nonvanishing scalar asymmetries  $A_d^{00+}(IM)$  and  $A_{ed}^{00+}(IM)$  marked by “√”.

Type	Current	00+	10+	11+	11-	20+	21+	21-	22+	22-
$A_d^{00+}(IM)$	PT-conserving	√				√	√		√	
	P-violating	√	√	√		√	√			
$A_{ed}^{00+}(IM)$	PT-conserving		√	√						
	P-violating	√	√	√		√	√		√	

We see that  $A_d^{00+}(00+)$ ,  $A_d^{00+}(20+)$ ,  $A_d^{00+}(21+)$ ,  $A_{ed}^{00+}(10+)$ , and  $A_{ed}^{00+}(11+)$  receive contributions from both parity violating as well as parity conserving pieces. It is therefore not possible to study any parity violating effects through observation of any of these asymmetry parameters as they will be completely swamped by the parity conserving contributions. The only way to study parity violating effects in e-d scattering with an unpolarized deuteron target will be through the asymmetry parameters  $A_d^{00+}(10+)$  and  $A_d^{00+}(11+)$  with unpolarized electron scattering and

through  $A_{ed}^{00+}(00+)$ ,  $A_{ed}^{00+}(20+)$ ,  $A_{ed}^{00+}(21+)$ , and  $A_{ed}^{00+}(22+)$  with polarized electron scattering. These are related to the vector polarization of the deuteron  $P_x$  and  $P_z$ , electron beam asymmetry  $\mathcal{A}$  and various tensor polarization asymmetries  $\mathcal{A}_{zz}$ ,  $\mathcal{A}_{xx/yy}$  and  $\mathcal{A}_{zx}$ . Observational quantities are defined in terms of the well known deuteron charge  $G_C(Q^2)$ , magnetic moment  $G_M(Q^2)$ , and quadrupole  $G_Q(Q^2)$  form factors as well as the new axial vector form factor  $G_A(Q^2)$  and electric form factor  $G_E(Q^2)$  which are defined in terms of various multipoles as follows:

$$G_C(Q^2) = \sqrt{\frac{4\pi}{3}} \frac{\beta}{1+\eta} C_0, \quad (104)$$

$$G_Q(Q^2) = \sqrt{\frac{3\pi}{2}} \frac{\beta}{\eta(1+\eta)} C_2, \quad (105)$$

$$G_{E/M}(Q^2) = \sqrt{\frac{\pi}{\eta(1+\eta)}} (E/M)_L, \quad (106)$$

$$G_A(Q^2) = \sqrt{\frac{\pi}{\eta(1+\eta)}} F_{E1}^A. \quad (107)$$

In the following we present results for these asymmetries and discuss the possibility to experimentally determine them.

### A. e-d Scattering with Unpolarized Electrons

In the case of elastic electron deuteron scattering with unpolarized electrons ( $h = 0$ ) (see Eqs. (46)-(48)) and unpolarized deuterons, the differential cross sections (Eq. (22)) and various tensor polarization components are dominated by the parity conserving contributions and there is no hope of observing any parity violating effects in these observables. However, the vector polarization  $P_x$  and  $P_z$  depend solely upon the parity violating contributions  $A_d^{00+}(10+)$  and  $A_d^{00+}(11+)$  (Eqs. (32)-(34)). Neglecting the parity violating contributions to the total cross section, we obtain

$$\begin{aligned} S_0 P_z &= \sqrt{\frac{2}{3}} S_0 A_d^{00+}(10+) \\ &= \frac{2}{3} \eta \left[ 1 + 2(1+\eta) \tan^2 \frac{\theta}{2} \right] [G_E + \tilde{G}_v G_A] G_M + \frac{4}{3} \sec \frac{\theta}{2} \tan \frac{\theta}{2} \eta \\ &\quad \times \sqrt{(1+\eta) \left( 1 + \eta \sin^2 \frac{\theta}{2} \right)} g_v^d \tilde{G}_a G_M^2, \end{aligned} \quad (108)$$

$$\begin{aligned} S_0 P_x &= -\frac{1}{\sqrt{3}} S_0 A_d^{00+}(11+) \\ &= -\frac{4}{3} \sec \frac{\theta}{2} \sqrt{\eta \left( 1 + \eta \sin^2 \frac{\theta}{2} \right)} [G_E + \tilde{G}_v G_A] [G_C + \frac{\eta}{3} G_Q] \\ &\quad - \frac{8}{3} \tan \frac{\theta}{2} \sqrt{\eta(1+\eta)} g_v^d \tilde{G}_a [G_C + \frac{\eta}{3} G_Q] G_M, \end{aligned} \quad (109)$$

where  $A_d^{00+}(11+)$  and  $A_d^{00+}(10+)$  are given in Eqs. (35) and (38). Both polarization observables  $P_x$  and  $P_z$  depend upon the contribution containing  $G_E + \tilde{G}_v G_A$  in addition to the Standard Model contribution proportional to  $\tilde{G}_a$  which is dominant. The contribution of  $G_E$  depends upon the p-wave component of the deuteron wave function while  $G_A$  depends upon the isoscalar axial vector form factor, which may be due to an isoscalar strangeness component of the nucleon or an anapole moment form factor, which both are expected to be small. We have considered these effects including a strangeness component in the magnetic moment form factor as well, which are parameterized as given in Eqs. (96), (97) and (101) with  $g_A^s = -0.19$ .

In Fig.2, we present the numerical values of  $P_x$  and  $P_z$  as a function of  $Q^2$  for various electron energies, corresponding to future studies of parity violating effects in electron deuteron scattering to be done at MAINZ [26] and JLAB [27]. We see that the effect of parity violation via P-odd admixtures is very small as compared to the contribution from  $\gamma$ -Z interference. In order to see the effect of the strangeness form factors in magnetic as well as axial vector form factors, we have considered three cases: (i)  $G_M^S \neq 0$ ,  $G_A^S = 0$ , and  $G_E^S = 0$ , (ii)  $G_M^S = 0$ ,  $G_A^S \neq 0$ , and  $G_E^S = 0$ , and (iii)  $G_M^S \neq 0$ ,  $G_A^S \neq 0$ , and  $G_E^S \neq 0$ , and show the effect in Figs. 3 and 4 as a function of  $Q^2$ , and as a function of  $\theta$ , for fixed  $Q^2$ , at various electron energies. For this purpose, we have neglected the  $G_E$  contribution to  $P_x$  and  $P_z$ , which is quite small. One readily notes that, if presently suggested values of  $G_M^S$  and  $G_A^S$  are taken, then the effect of a

strangeness component in the axial vector and magnetic moments are opposite in sign and tend to cancel each other. Moreover, the magnetic effect of a strangeness component in the magnetic moment is larger than the effect of including a strangeness component in the axial vector form factor if presently suggested values for their magnitudes are used. The parity violating effects in elastic deuteron scattering have been considered before by many authors [33]-[37]. But none of them has calculated the vector polarization of the recoil deuteron, although it is discussed qualitatively by Ramachandran and Singh [35].

### B. e-d Scattering with Polarized Electrons

In the case of elastic electron scattering with polarized electrons ( $h \neq 0$ ) on unpolarized deuterons, the vector polarization components of the recoil deuteron are dominated by the parity conserving contributions  $A_{ed}^{00+}(11+)$  and  $A_{ed}^{00+}(10+)$ , and thus no useful information on parity violation can be obtained from observing  $P_x$  and  $P_z$ . In this case the purely parity violating contributions  $A_{ed}^{00+}(00+)$ ,  $A_{ed}^{00+}(20+)$ ,  $A_{ed}^{00+}(21+)$ , and  $A_{ed}^{00+}(22+)$  (see Table III) can be determined through appropriately defined asymmetries of the differential cross section and the recoil tensor polarizations. For example, defining the beam asymmetry as

$$\mathcal{A} = \frac{1}{h} \frac{\left[ \left( \frac{d\sigma}{d\Omega} \right)_+ - \left( \frac{d\sigma}{d\Omega} \right)_- \right]}{\left[ \left( \frac{d\sigma}{d\Omega} \right)_+ + \left( \frac{d\sigma}{d\Omega} \right)_- \right]}, \quad (110)$$

where  $\left( \frac{d\sigma}{d\Omega} \right)_\pm$  denotes the differential cross section with right handed ( $h = +1$ ) and left handed ( $h = -1$ ) polarized electrons, then one obtains  $\mathcal{A} = A_{ed}^{00+}(00+)$ . It is given, separating it into  $\mathcal{A}_Z$  and  $\mathcal{A}_\gamma$ , by

$$\mathcal{A} = \mathcal{A}_Z + \mathcal{A}_\gamma = A_{ed}^{00+}(00+) = 2g_v^d \tilde{G}_a + \frac{8}{3S_0} \sec \frac{\theta}{2} \tan \frac{\theta}{2} \eta \sqrt{(1+\eta) \left( 1 + \eta \sin^2 \frac{\theta}{2} \right)} \left( G_E + \tilde{G}_v G_A \right) G_M. \quad (111)$$

In Figs. 5 and 6, we show this asymmetry as a function of  $Q^2$  as well as a function of the scattering angle  $\theta$ , for fixed  $Q^2$ , at various energies available for experiments at MAINZ [26] and JLAB [27]. We have also studied the effect of a strangeness form factor on these asymmetries for various cases discussed in subsection V A using the functional form of the strangeness form factor given in Eqs. (96), (97) and (101). For this purpose, we have again neglected the contribution of  $G_E$ . We see from Eq. (111) that apart from the dominant contribution proportional to  $\tilde{G}_a$ , the additional contribution depends upon  $G_E + \tilde{G}_v G_A$ , the same combination of weak form factors as encountered in  $P_x$  and  $P_z$ , but with a different angular dependence.

We see that the dominant contribution to  $\mathcal{A}$  comes from the  $\gamma$ -Z interference whereas the contribution of  $G_E$ , coming from the parity violating effect in the wave function, is small. This is similar to the results found in quasi-elastic electron-deuteron scattering. The axial vector contribution comes mainly from the strangeness form factors, as the radiative correction to the isoscalar piece is quite small unlike the isovector case. Therefore, the observation of  $\mathcal{A}$  in elastic polarized electron deuteron scattering is best suited to study the axial vector strangeness form factor, if one has knowledge of the strangeness in the magnetic form factor. In Figs. 5 and 6, we also show the effect of a nonzero strangeness contribution in the magnetic and axial vector form factors. We see that with the presently suggested value of  $G_M^S$  and  $G_A^S$ , the effect on  $\mathcal{A}$  from  $G_M^S$  and  $G_A^S$  are opposite in nature. The strangeness contribution to  $\mathcal{A}$  is dominated by  $G_M^S$  as compared to  $G_A^S$ . This asymmetry was discussed earlier by many authors [33, 34, 36, 37]. While Porrman [33] discusses  $\mathcal{A}_z$  and  $\mathcal{A}_\gamma$ , he does not take into account the effect of a strangeness content in the vector and axial vector currents, the work of Hwang and Henley [28] discusses only  $\mathcal{A}_z$ . On the other hand, the work of Pollock [34] and Frederico et al. [36] considers the asymmetry  $\mathcal{A}_z$  by taking into account an isoscalar axial vector contribution, while the effect of  $\mathcal{A}_\gamma$  is neglected. Thus the present study is the first one, which takes into account all contributions and presents results for  $\mathcal{A}_z$  and  $\mathcal{A}_\gamma$  for the electron energies of current interest.

The parity violation effects also enter in the tensor polarizations of the recoil deuteron through  $A_{ed}^{00+}(20+)$ ,

$A_{ed}^{00+}(21+)$ , and  $A_{ed}^{00+}(22+)$  which are given in terms of Sachs form factors

$$\begin{aligned} S_0 A_{ed}^{00+}(20+) = & -\frac{8\sqrt{2}}{3} g_v^d \tilde{G}_a \left[ G_C + \frac{\eta}{3} G_Q \right] G_Q - \frac{\sqrt{2}}{3} \eta \left[ 1 + 2(1+\eta) \tan^2 \frac{\theta}{2} \right] g_v^d \tilde{G}_a G_M^2 \\ & - \frac{2\sqrt{2}}{3} \eta \sec \frac{\theta}{2} \tan \frac{\theta}{2} \sqrt{(1+\eta) \left( 1 + \eta \sin^2 \frac{\theta}{2} \right)} \left[ G_E + \tilde{G}_v G_A \right] G_M, \end{aligned} \quad (112)$$

$$\begin{aligned} S_0 A_{ed}^{00+}(21+) = & \frac{8}{\sqrt{3}} \eta \sec \frac{\theta}{2} \sqrt{\eta \left( 1 + \eta \sin^2 \frac{\theta}{2} \right)} g_v^d \tilde{G}_a G_M G_Q \\ & + \frac{4}{\sqrt{3}} \eta \tan \frac{\theta}{2} \sqrt{\eta(1+\eta)} \left[ G_E + \tilde{G}_v G_A \right] G_Q, \end{aligned} \quad (113)$$

$$S_0 A_{ed}^{00+}(22+) = -\frac{2}{\sqrt{3}} \eta g_v^d \tilde{G}_a G_M^2. \quad (114)$$

While  $A_{ed}^{00+}(22+)$  is given solely in terms of the Standard Model parameters (see Eq. (114)) the other components involve the isoscalar axial vector piece and the parity violating hadronic interaction in the combination  $G_E + \tilde{G}_v G_A$ , the same which occurs in the asymmetry  $\mathcal{A}$ . In order to extract these one has to measure tensor polarization asymmetries. For example, one may define the tensor polarization asymmetry  $\mathcal{A}_{zz}^p$  by

$$\mathcal{A}_{zz}^p = \frac{1}{h} \frac{p_{zz}(\uparrow) - p_{zz}(\downarrow)}{p_{zz}(\uparrow) + p_{zz}(\downarrow)} = \frac{A_{ed}^{00+}(20+)_{pv}}{A_d^{00+}(20+)_{pc+pv}}. \quad (115)$$

In Fig. 7, we show the results for  $\mathcal{A}_{zz}^p$  as a function of  $Q^2$  for various values of electron energy  $E = 125, 200, 315, 361$ , and 687 MeV. We also show the effect of including strangeness form factors in the magnetic as well as axial vector form factors. We note that  $\mathcal{A}_{zz}^p$  remains constant for a large range of  $Q^2$  values. This is not surprising since  $\mathcal{A}_{zz}^p$  is a ratio of  $A_{ed}^{00+}(20+)$  and  $A_d^{00+}(20+)$ , which have a similar  $Q^2$  dependence, though differing in magnitude by a large amount.

## VI. SUMMARY AND CONCLUSIONS

Parity violating electron scattering experiments are being done by the SAMPLE, HAPPEX, G0 and A4 collaborations at MIT, JLAB and MAINZ from proton, deuteron and  ${}^4\text{He}$  targets in the energy region of a few hundred MeV. These experiments are expected to answer questions about the strange form factors of the nucleon and radiative corrections to the axial vector couplings which may also lead to an understanding of anapole moments.

In this paper we have examined parity violating observables in the scattering of unpolarized and polarized electrons from unpolarized deuterons. The parity violating observables receive contributions from the  $\gamma$ -Z interference in the Standard Model and also from the parity violating electromagnetic coupling which are induced by the parity violating components in the deuteron wave function. In almost all parity violating observables, the contribution of the parity violating electromagnetic coupling is found to be small compared to the  $\gamma$ -Z interference contribution. In the case of unpolarized electron scattering, the nonvanishing components of the parity violating recoil vector polarization of the deuteron have been studied at electron energies  $E = 125, 200, 315, 362$ , and 687 MeV, relevant for the ongoing experiments as a function of  $Q^2$ . In addition, for a fixed value of  $Q^2$ , these polarizations have been studied as a function of the scattering angle. The effects of a nonzero strangeness component in the magnetic form factor as well as the axial vector form factor have been calculated. These effects are found to be important in backward direction. In the case of the axial vector coupling, the radiative corrections are known to be large in the isovector component and play an important role in the analysis of quasi-elastic electron deuteron scattering. For elastic electron deuteron scattering, where isoscalar terms contribute, these radiative corrections are small, thus making it a suitable method to study the strangeness form factor.

In the case of polarized electron scattering from unpolarized deuterons, the parity violating electron beam asymmetry and the tensor polarization asymmetries of the recoiling deuteron have been studied. For the electron beam asymmetry, the results are presented as a function of  $Q^2$  for various energies. Furthermore, the asymmetries have been studied as a function of the scattering angle at fixed  $Q^2$ . The effect of including a nonzero component in the magnetic form factor  $G_M^s(Q^2)$  as well as an axial form factor  $G_A^s(Q^2)$  have been investigated too.

In all these cases, the contributions to the parity violating observables are dominated by the  $\gamma$ -Z interference term of the Standard Model where the leptonic axial vector current interacts with the hadronic isoscalar axial current. There

are additional contributions, though small, coming from the parity violating electromagnetic coupling as well as from the axial vector isoscalar hadronic current which interacts with the leptonic vector current. This isoscalar hadronic current may be due to the strangeness component of the nucleon in the Standard Model or due to the structure of weak interactions beyond the Standard Model. It is found that the parity violating observables are important in the backward direction. The presented results for various parity violating asymmetries might be helpful in analysing future experiments on electron-deuteron scattering being done at MAINZ, MIT and JLAB.

## VII. ACKNOWLEDGMENT

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- [1] D. H. Beck, Phys. Rev. D **39**, 3248 (1989).
- [2] R. D. McKeown, Phys. Lett. B **219**, 140 (1989).
- [3] D. Kaplan and A. Manohar, Nucl. Phys. B **310**, 527 (1988).
- [4] J. Ashman et al., Phys. Lett. B **206**, 364 (1988).
- [5] L. Ahrens et al., Phys. Rev. D **35**, 785 (1987).
- [6] J. F. Donoghue and C.R. Nappi, Phys. Lett. B **168**, 105 (1986).
- [7] E. J. Beise, M.L. Pitt and D.T. Spayde, Prog. Part. Nucl. Phys. **54**, 289 (2005).
- [8] D. T. Spayde et al., Phys. Rev. Lett. **84**, 1106 (2000).
- [9] K. A. Aniol et al., Phys. Rev. C **69**, 065501 (2004).
- [10] K. A. Aniol et al., Phys. Lett. B **635**, 275 (2006).
- [11] A. Acha et al., Phys. Rev. Lett. **98**, 032301 (2007).
- [12] F. E. Mass et al., Phys. Rev. Lett. **93**, 022002 (2004).
- [13] F. E. Mass et al., Phys. Rev. Lett. **94**, 152001 (2005).
- [14] D. S. Armstrong et al., Phys. Rev. Lett. **95**, 092001 (2005).
- [15] T. M. Ito et al., Phys. Rev. Lett. **92**, 102003 (2004).
- [16] K. A. Aniol et al., Phys. Rev. Lett. **96**, 022003 (2006).
- [17] J. Liu, R.D. McKeown and M.J. Ramsey-Musolf, Phys. Rev. C **76**, 025202 (2007).
- [18] R. D. Young, J. Roche, R. D. Carlini, and A. W. Thomas, Phys. Rev. Lett. **97**, 102002 (2006).
- [19] S. F. Pate, Phys. Rev. Lett. **92**, 082002 (2004).
- [20] M. J. Musolf, T. W. Donniley, J. Dubach, S. J. Pollock, S. Kowalski, and E. J. Beise, Phys. Rep. **239**, 1 (1994).
- [21] S.-L. Zhu, S. J. Puglia, B. R. Holstein, M. J. Ramsey-Musolf, Phys. Rev. D **62**, 033008 (2000).
- [22] W. C. Haxton, E. M. Henley, M. J. Musolf, Phys. Rev. Lett. **63**, 949 (1989).
- [23] M. J. Musolf and B. R. Holstein, Phys. lett. B **242**, 461 (1990).
- [24] R. Schiavilla, J. Carlson, and M. Paris, Phys. Rev. C **67**, 032501(R) (2003).
- [25] C.P. Liu, G. Prezeau, and M.J. Ramsey-Musolf, Phys. Rev. C **67**, 035501 (2003).
- [26] S. Baunack, Eur. Phys. J. A **32**, 457 (2007).
- [27] J. Roche, contribution presented at *Gordon Conference on Photonuclear Reactions* (not published) (2006).
- [28] W. Y. P. Hwang, E. Henley and G. Miller, Ann. Phys. **137**, 378 (1981).
- [29] E. Hadjimichael, G. I. Poulis, and T. W. Donnelly, Phys. Rev. C **45**, 2666 (1992).
- [30] B. Mosconi and P. Ricci, Phys. Rev. C **55**, 3115 (1997).
- [31] G. Küster and H. Arenhövel, Nucl. Phys. A **626**, 911 (1997).
- [32] R. Schiavilla, J. Carlson, and M. Paris, Phys. Rev. C **70**, 044007 (2004).
- [33] M. Porrmann, Nucl. Phys. A **360**, 251 (1981).
- [34] S. J. Pollock, Phys. Rev. D **42**, 3010 (1990).
- [35] G. Ramachandran and S.K. Singh, Phys. Rev. D **18**, 1441 (1978).
- [36] T. Frederico, E.M. Henley, and G.A. Miller, Nucl. Phys. A **533**, 617 (1991).
- [37] W. Y. P. Hwang and E. M. Henley, Ann. Phys. **129**, 47 (1980).
- [38] H. Arenhövel and S.K. Singh, Eur. Phys. J. A **10**, 183 (2001).
- [39] R. Machleidt, K. Holinde and C. H. Elster, Phys. Rep. **149**, 1 (1987).
- [40] B. Desplanques, J.F. Donoghue, and B.R. Holstein, Ann. Phys. (N.Y.) **124**, 449 (1980).
- [41] V. M. Dubovik and S. Z. Zenkin, Ann. Phys. (NY), **172**, 100 (1986).
- [42] G. B. Feldman, G. A. Crawford, J. Dubach and B. R. Holstein, Phys. Rev. **C43**, 863 (1991).
- [43] H. Arenhövel, M. Danos and H. T. Williams, Nucl. Phys. A **162**, 12 (1971).
- [44] J. D. Walecka, *Theoretical Nuclear and Subnuclear Physics*, Oxford University Press, USA, (1995).

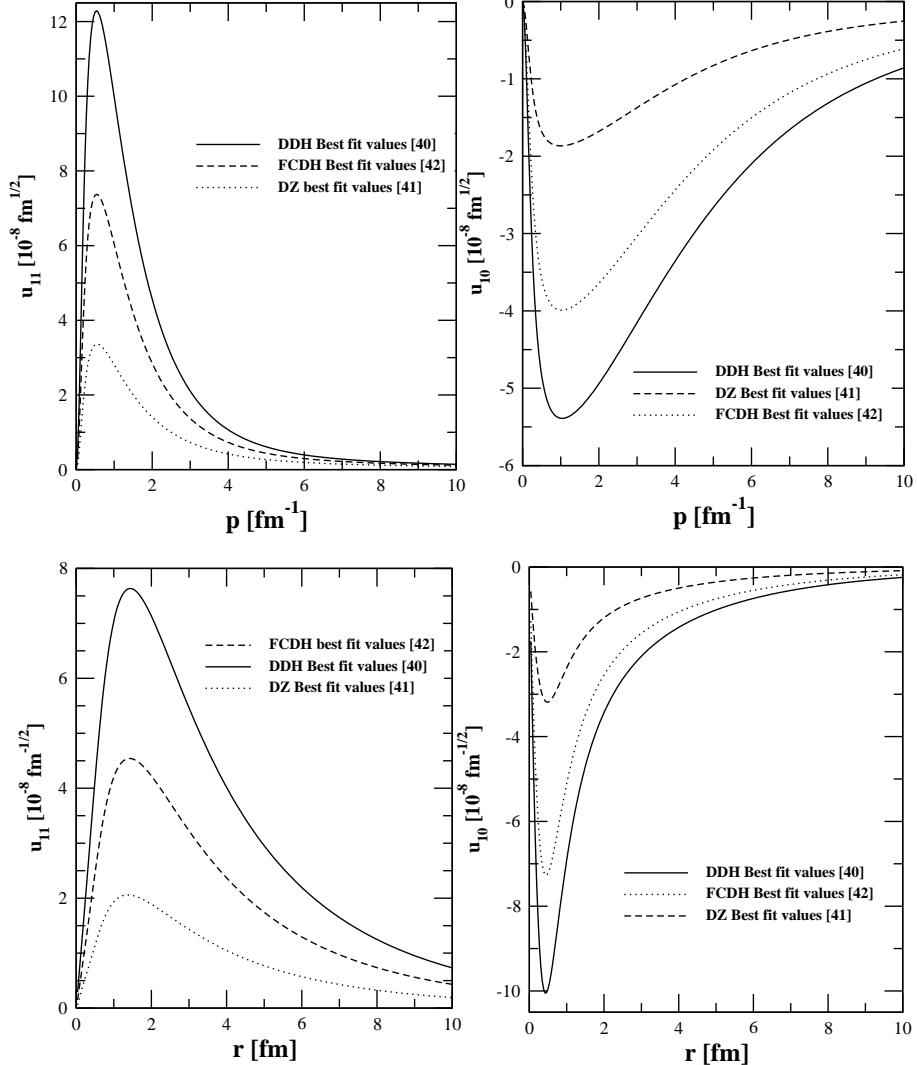


FIG. 1: Radial part of p-wave components of the deuteron wave function in momentum space and in configuration space. Separately shown are the sensitivity to the parameters of the weak nucleon-nucleon potential.

- [45] W. M. Alberico, S. M. Bilenky and C. Maieron, Phys. Rep. **358**, 227 (2002).
- [46] S. K. Singh, Nucl. Phys. B (Proc. Suppl.) **112**, 77 (2002).
- [47] D. H. Beck and B. R. Holstein, Int. J. Mod. Phys. **E10**, 1 (2001).
- [48] H. -W. Hammer, S. J. Puglia, M. J. Ramsey-Musolf and S. -L. Zhu, Phys. Lett. B **562**, 208 (2003).
- [49] D. Leinweber and A. W. Thomas, Phys. Rev. D **62**, 074505 (2000).
- [50] R. Lewis, W. Wilcox and R. M. Woloshyn, Phys. Rev. D **67**, 013003 (2003).

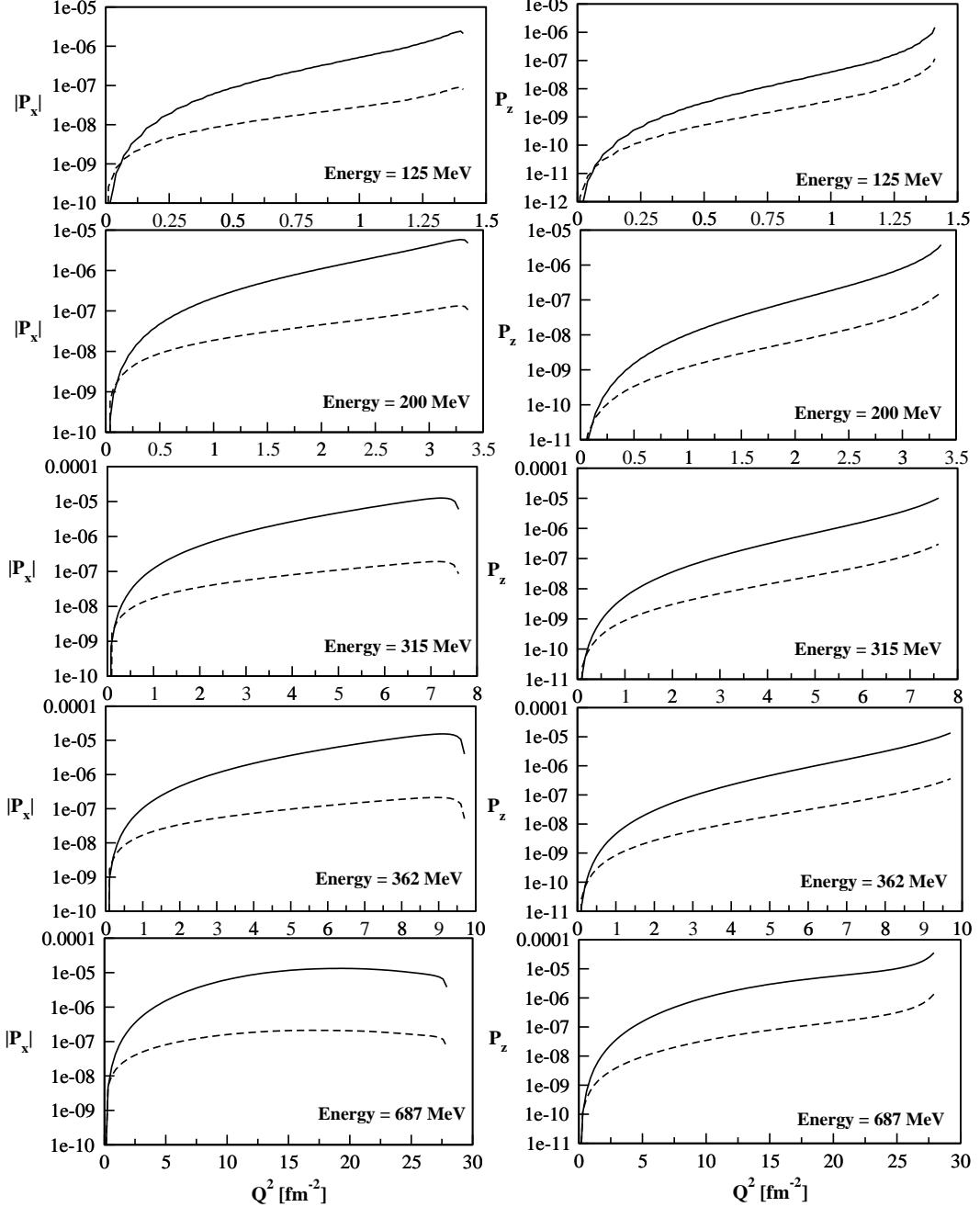


FIG. 2: Recoil vector polarizations  $P_x$  (left panel) and  $P_z$  (right panel) as a function of  $Q^2$  for various energies  $E = 125, 200, 315, 361$ , and  $687$  MeV. The solid curves represent the contribution of Z-exchange and the dashed curves the contribution of the  $\gamma$ -Z interference due to P state admixtures in the deuteron wave function.

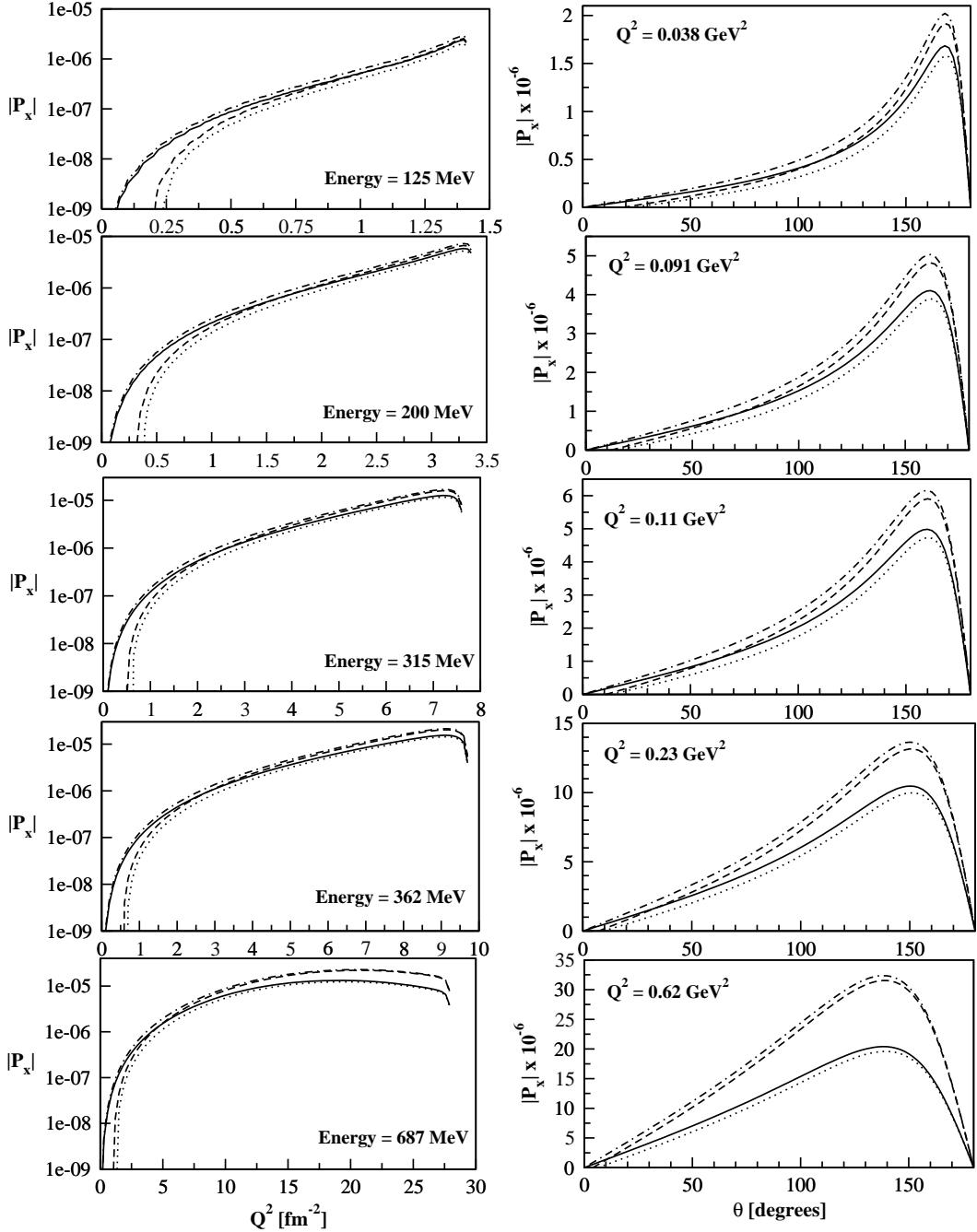


FIG. 3: Recoil vector polarization  $P_x$  by Z-exchange as a function of  $Q^2$  for various energies  $E = 125, 200, 315, 361$ , and  $687$  MeV (left panel), and as a function of  $\theta$  for fixed  $Q^2 = 0.038, 0.091, 0.11, 0.23$ , and  $0.62$   $GeV^2$ . The effect of strangeness form factors in magnetic as well as axial vector form factors are shown for three cases: (i)  $G_M^s \neq 0$  and  $G_A^s = 0$  (dashed-dotted), (ii)  $G_M^s = 0$  and  $G_A^s \neq 0$  (dotted), and (iii)  $G_M^s \neq 0$  and  $G_A^s \neq 0$  (dashed). The solid curves represent the case when  $G_M^s = 0$  and  $G_A^s = 0$ .

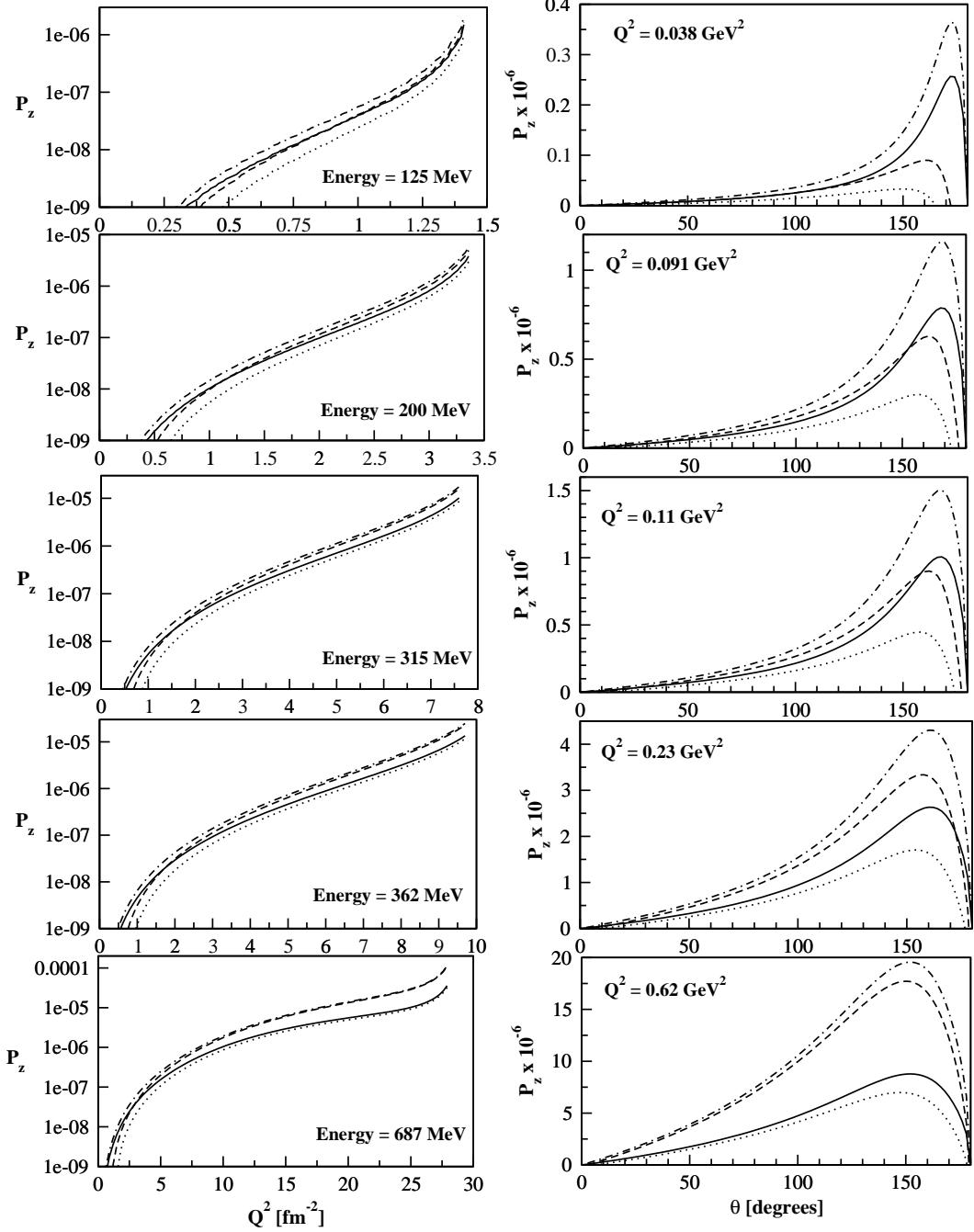


FIG. 4: Recoil vector polarizations  $P_z$  by Z-exchange as a function of  $Q^2$  for various energies  $E = 125, 200, 315, 361$ , and  $687$  MeV (left panel), and as a function of  $\theta$  for fixed  $Q^2 = 0.038, 0.091, 0.11, 0.23$ , and  $0.62$   $\text{GeV}^2$ . The effect of strangeness form factors in magnetic as well as axial vector form factors are shown for three cases: (i)  $G_M^s \neq 0$  and  $G_A^s = 0$  (dashed-dotted), (ii)  $G_M^s = 0$  and  $G_A^s \neq 0$  (dotted), and (iii)  $G_M^s \neq 0$  and  $G_A^s \neq 0$  (dashed). The solid curves represent the case when  $G_M^s = 0$  and  $G_A^s = 0$ .

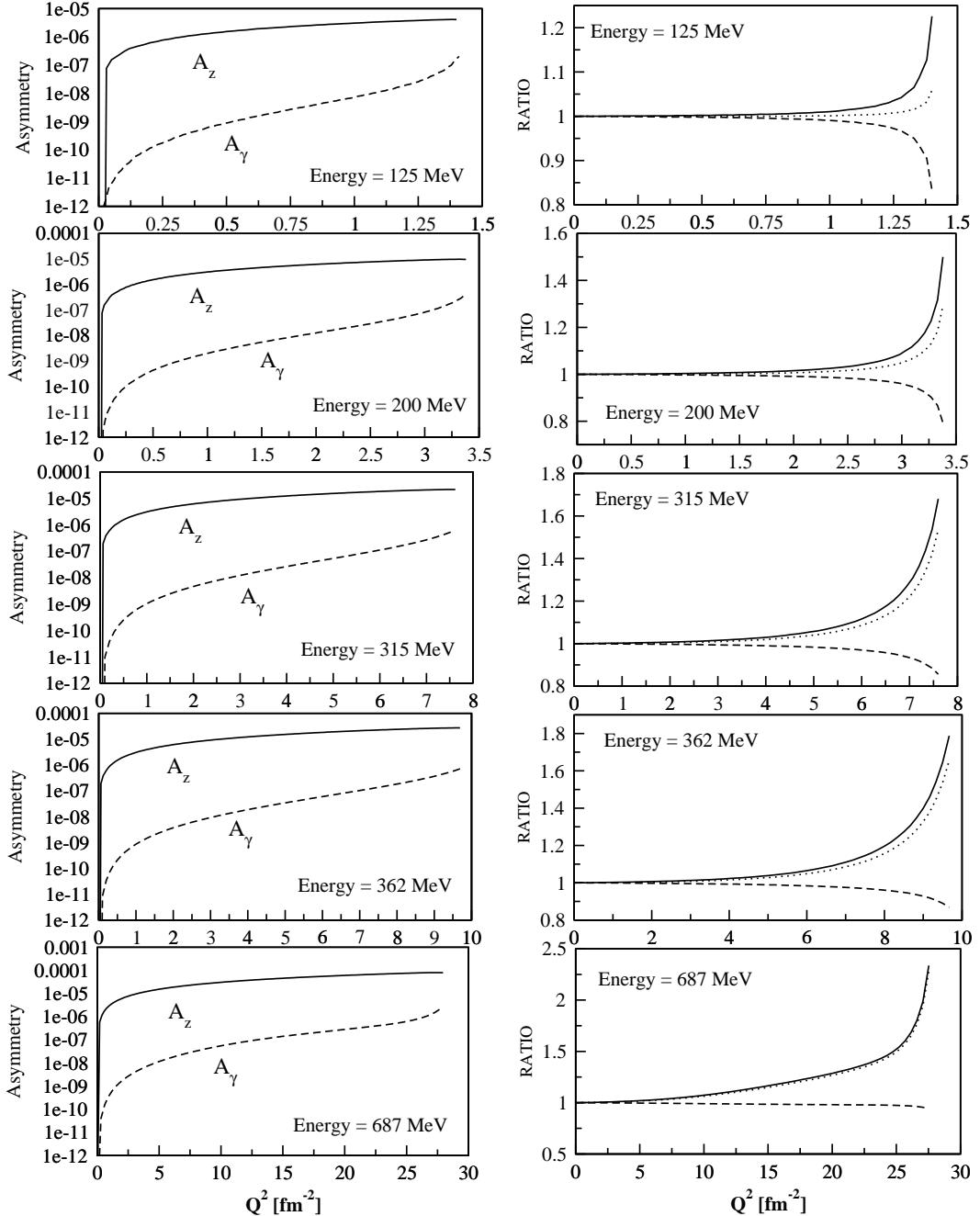


FIG. 5: Elastic deuteron asymmetry by Z [ $A_z$  (solid)] and photon exchange [ $A_\gamma$  (dashed)] as a function of  $Q^2$  for various incident energies  $E = 125, 200, 315, 361$ , and  $687$  MeV (left panel). The effect of strangeness form factors in magnetic  $G_M^s$  as well as axial vector  $G_A^s$  form factor are shown as a ratio of strange contributions (right panel): (i)  $G_M^s \neq 0$  and  $G_A^s = 0$  (solid), (ii)  $G_M^s = 0$  and  $G_A^s \neq 0$  (dashed), and (iii)  $G_M^s \neq 0$  and  $G_A^s \neq 0$  (dotted), with no strange contributions  $G_M^s = 0$  and  $G_A^s = 0$ .

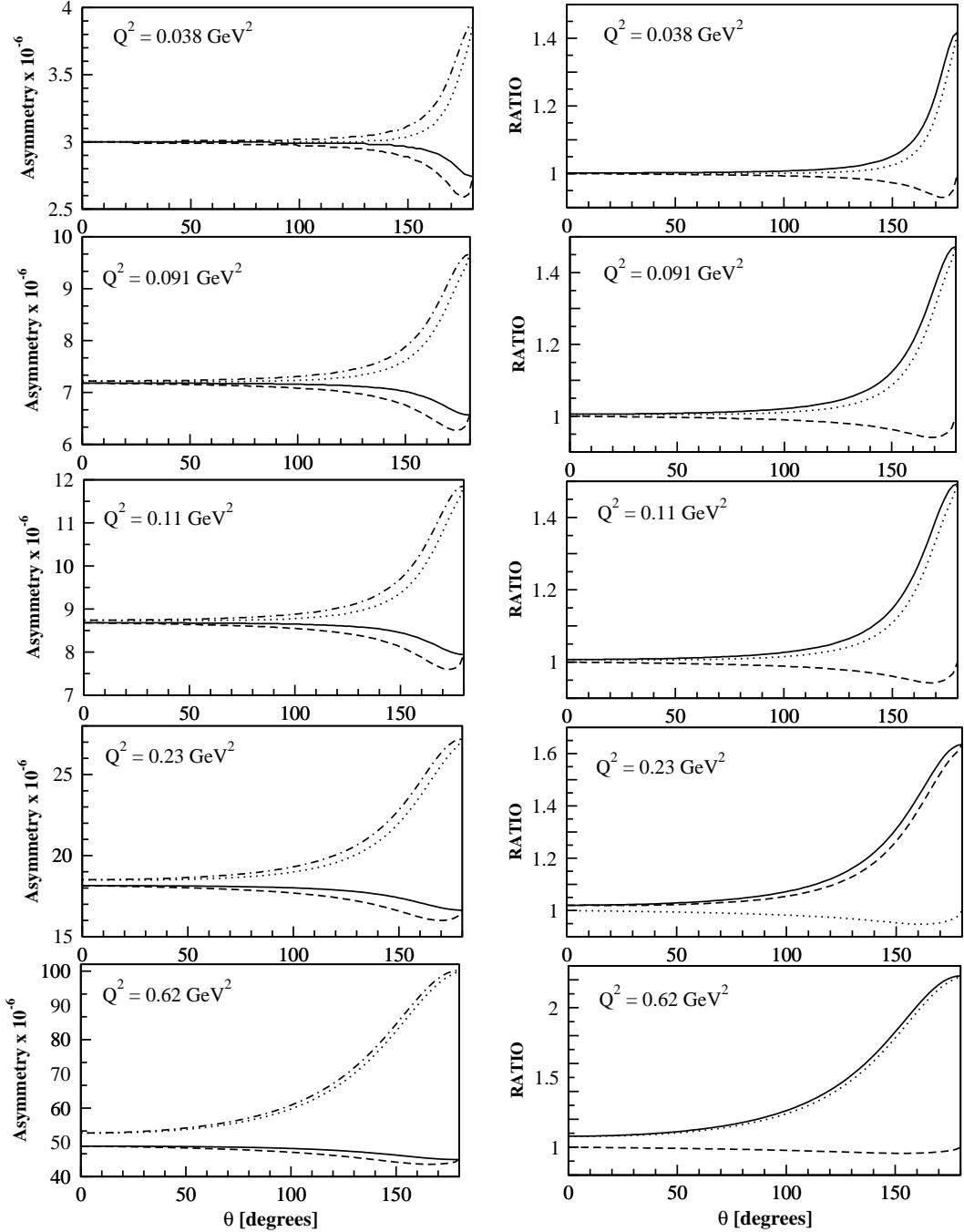


FIG. 6: Elastic deuteron asymmetry by Z-exchange [ $\mathcal{A}_z$  (solid)] as a function of  $\theta$  for fixed  $Q^2 = 0.038, 0.091, 0.11, 0.23$ , and  $0.62 \text{ GeV}^2$  (left panel). The effect of strangeness form factors in magnetic as well as axial vector form factors have been shown for three cases: (i)  $G_M^s \neq 0$  and  $G_A^s = 0$  (dashed-dotted), (ii)  $G_M^s = 0$  and  $G_A^s \neq 0$  (dashed), and (iii)  $G_M^s \neq 0$  and  $G_A^s \neq 0$  (dotted). The solid curves represent the case when  $G_M^s = 0$  and  $G_A^s = 0$ . In the right panel of the figure the effect of strangeness form factors have been shown as a ratio of strange contributions (i)  $G_M^s \neq 0$  and  $G_A^s = 0$  (solid), (ii)  $G_M^s = 0$  and  $G_A^s \neq 0$  (dashed), and (iii)  $G_M^s \neq 0$  and  $G_A^s \neq 0$  (dotted), with no strange contributions  $G_M^s = 0$  and  $G_A^s = 0$ .

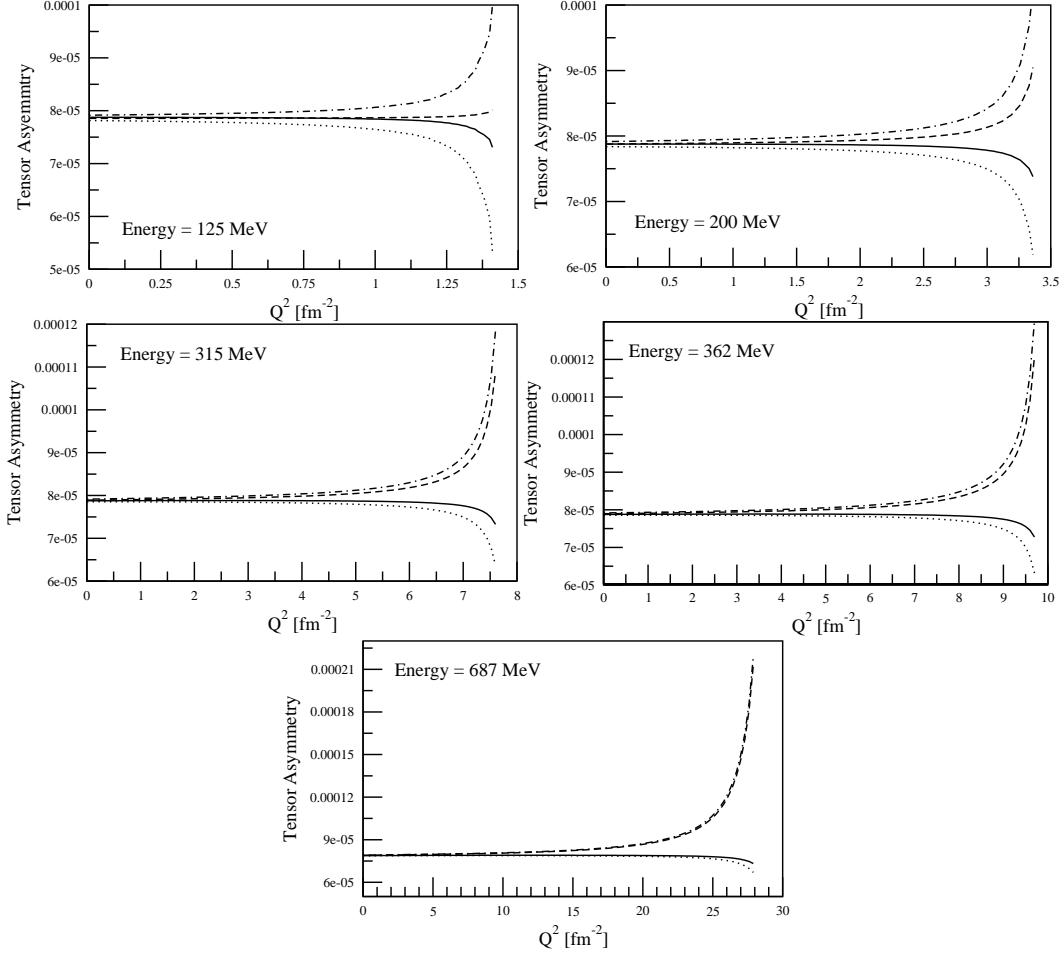


FIG. 7: Recoil tensor asymmetry  $A_{zz}^p$  from Z-exchange as a function of  $Q^2$  for various incident energies  $E = 125, 200, 315, 361$ , and  $687$  MeV. The effect of strangeness form factors in magnetic as well as axial vector form factors are shown for three cases: (i)  $G_M^s \neq 0$  and  $G_A^s = 0$  (dashed-dotted), (ii)  $G_M^s = 0$  and  $G_A^s \neq 0$  (dotted), and (iii)  $G_M^s \neq 0$  and  $G_A^s \neq 0$  (dashed). The solid curves represent the case when  $G_M^s = 0$  and  $G_A^s = 0$ .